

# Summer of Science

Quantum Computing & Information - P02

Shaik Awez Mehtab

**Roll No:** 23b1080

**Mentor:** Kanishk Modi

## Final Report



IIT Bombay

July 27, 2024

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## **Part I**

### **Week 1**

# Chapter 1

## Linear Algebra

Usual Linear algebra notation is changed to **Dirac Notation**.

We consider vector spaces in  $\mathbb{C}^n$ , with additive & multiplicative properties on  $\mathbb{C}^n$  too, i.e  $\mathbb{C}^n$  is  $n$  dimensional. Here, the scalars are **complex**.

### 1.1 Dirac Notation

A vector  $\psi$  is denoted by  $|\psi\rangle$ <sup>1</sup>.  $|\psi\rangle$  is referred to as ket, it's vector dual  $\langle\psi|$  is referred to as bra.  $|\psi\rangle$  is nothing but an  $n$ -tuple of complex numbers,  $(\psi_1, \dots, \psi_n)$  or a column vector

$$\begin{bmatrix} \psi_1 \\ \vdots \\ \psi_n \end{bmatrix}$$

few more dirac notations:

Notation	Description
$z^*$	Complex conjugate of $z$ , $z \in \mathbb{C}$ . eg. $(1 + i)^* = 1 - i$
$\langle\phi \psi\rangle$	Inner product (Dot product) of $ \phi\rangle$ and $ \psi\rangle$
$ \psi\rangle \otimes  \phi\rangle$ or $ \psi\rangle \langle\phi $	Outer product of $ \psi\rangle$ and $ \phi\rangle$
$A^\dagger$	Hermitian conjugate or Adjoint of $A$ , $= (A^T)^*$
$\langle\psi A \phi\rangle$	Inner product between $\psi$ and $A \phi\rangle$ or, Inner product between $ \phi\rangle$ and $A^\dagger \psi\rangle$

Table 1.1: Useful *Dirac* notations.

Revise spanning set, span, linear independence, basis, dimension. They're similar to the case of real scalars.

I'll write down few important definitions from linear algebra (and few facts too):

### 1.2 Linear Operators

A function  $A$ , from vector space  $V$  to  $W$ , which obey

$$A\left(\sum_i a_i |v_i\rangle\right) = \sum_i a_i A(|v_i\rangle)$$

---

<sup>1</sup> $|0\rangle$  isn't the zero vector, zero vector is denoted by just 0

in simpler words, operator on a linear combination of inputs is same linear combination of operator on each input

When based of inputs & output are specified, using coordinate vectors, A linear operator is equivalent to matrix multiplication.

### 1.2.1 Pauli Matrices

These are 4 fundamentally important matrices in quantum information & computation.

$$\begin{aligned}\sigma_0 = I &= \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} & \sigma_1 = \sigma_x = X &= \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix} \\ \sigma_2 = \sigma_y = Y &= \begin{bmatrix} 0 & -i \\ i & 0 \end{bmatrix} & \sigma_3 = \sigma_z = Z &= \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix}\end{aligned}$$

## 1.3 Inner Products

This is a bit different from it's real counterpart, due to the reason below.

$$\langle v|w \rangle = (|v\rangle, |w\rangle) = \sum_i v_i^* w_i = [v_1^* \cdots v_n^*] \begin{bmatrix} w_1 \\ \vdots \\ w_n \end{bmatrix}$$

the conjugate comes from the fact that this should denoted length squared (for equal vectors). the formal reason is:

Inner product between two vectors  $|v\rangle$  and  $|w\rangle \in V$  is a function  $(\cdot, \cdot) : VV \rightarrow \mathbb{C}$  which satisfies:

1.  $(|v\rangle, \sum_i \lambda_i |w_i\rangle) = \sum_i \lambda_i (|v\rangle, |w_i\rangle)$
2.  $(|v\rangle, |w\rangle) = (|w\rangle, |v\rangle)^*$
3.  $(|v\rangle, |v\rangle) \geq 0, 0 \text{ only when } |v\rangle = 0$

$$(\sum_i \lambda_i w_i, v) = \sum_i \lambda_i^* (w_i, v) \tag{1.1}$$

Any vector space that has a defined inner product is called an **inner product space** or **Hilbert space** (if finite dimension) in QIC.

Revise orthogonality, norm, orthonormal set, gram-schmidt process. they're same as in real case. We'll consider only orthonormal based for linear operators, and assume same spaces have same bases.

### 1.3.1 Dual of a vector

it would've been evident from dot product that  $\langle v| = (|v\rangle^T)^*$  i.e if  $|v\rangle = (v_1, \dots, v_n)^T$ , then

$$\langle v| = [v_1^* \cdots v_n^*]$$

## 1.4 Outer Product

outer product of  $|v\rangle \in V$  &  $|w\rangle \in W$  is  $|v\rangle \langle w|$

it can be seen that this is a matrix, hence it's a linear operator from  $W$  to  $V$ . let  $|w'\rangle \in W$ ,

$$\underbrace{|v\rangle \langle w|}_{\text{operator}} |w'\rangle = |v\rangle \underbrace{\langle w|w'\rangle}_{\text{scalar}}$$

### 1.4.1 Completeness Relation

Continuing from yesterday, there's this not so trivial property using outer product, called **completeness relation**

$$\sum_i |i\rangle \langle i| = I$$

Here,  $|i\rangle$ 's form an orthonormal basis, they can be anything. one nice use of this is, if we have vector spaces  $V$  &  $W$ , which have bases  $|v_1\rangle, \dots, |v_m\rangle$  and  $|w_1\rangle, \dots, |w_n\rangle$  then any linear transform  $A : V \rightarrow W$  which is of size  $nm$  can be written as a linear combination of

$$|w_1\rangle \langle v_1|, \dots, |w_n\rangle \langle v_1|, \dots, |w_n\rangle \langle v_m|$$

each of which are  $nm$ . The linear combination is given by

$$A = \sum_{ij} \underbrace{\langle w_j | A | v_i \rangle}_{\text{scalar}} \underbrace{|w_j\rangle \langle v_i|}_{\text{one of the basis matrix}}$$

eg. I can write the pauli matrix  $\sigma_0$  as

$$1 * |0\rangle \langle 0| + 0 * |0\rangle \langle 1| + 0 * |1\rangle \langle 0| + 1 * |1\rangle \langle 1|$$

### 1.4.2 Cauchy Schwartz Inequality

For any two vectors  $|v\rangle, |w\rangle$  of same dimension, the following inequality is always true

$$\langle v|v\rangle \langle w|w \geq \langle v|w\rangle \langle w|v \rangle = |\langle v|w\rangle|^2$$

here's a short proof I've written.  $|0\rangle, \dots, |n\rangle$  are orthonormal based I've considered/can be constructed by gram schmidt process.

## 1.5 Eigenvectors & Eigenvalues

it's the same as you know, just a slight change in notation. we represent both eigenvalue and eigenvector label as  $v$ .

$$\underbrace{A}_{\text{linear operator}} |v\rangle = \underbrace{v}_{\text{scalar}} |v\rangle$$

and the characteristic function is  $c(\lambda) = \det(A - \lambda I)$ . it depends only on the operator, not on how it's matrix is represented.

Eigenspace of an eigenvalue  $v$  is the set of vectors which have eigenvalue  $v$ . read diagonalisability, orthonormal decomposition, degenerate eigenvectors.

## 1.6 Special Operators

### 1.6.1 Adjoint & Hermitian operators

for any linear operator  $A$  on a vector space  $V$ , there exists a unique linear operator  $A^\dagger$  such that,  $\forall |v\rangle, |w\rangle \in V$

$$(|v\rangle, A|w\rangle) = (A^\dagger|v\rangle, |w\rangle)$$

I've deduced that  $A^\dagger = (A^T)^*$  here, when  $A$  is represented as a matrix. This is actually same as dual of a vector. Few properties:

1.  $(AB)^\dagger = B^\dagger A^\dagger$
2.  $|v\rangle^\dagger = \langle v|$

Thus you can see  $(A|v\rangle)^\dagger = \langle v|A^\dagger$ , hence the above thing is evident.

It is **anti-linear**.

$$\left( \sum_i a_i A_i \right)^\dagger = \sum_i a_i^* A_i^\dagger$$

A **Hermitian** or a **self-adjoint operator** is a linear operator such that  $A^\dagger = A$ .

### 1.6.2 Projectors

these are nice. suppose  $W$  is a  $k$ -dimensional subspace of  $V$ , which is  $d$ -dimensional, we can always construct, using *gram-schmidt* process, an orthonormal basis  $|1\rangle, \dots, |d\rangle$  of  $V$  such that  $|1\rangle, \dots, |k\rangle$  is an orthonormal basis of  $W$ . Now we can project any vector in  $V$  onto  $W$ , using this operator:

$$P = \sum_{i=1}^k |i\rangle \langle i|$$

cross check this. This operator  $P$  is called the **projector**. It is also Hermitian. and we'll use  $P$  to refer to  $V$ .  $Q = I - P$  is the **orthogonal complement** of  $P$ .

### 1.6.3 Normal Operators

An operator  $A$  is normal if  $A^\dagger A = AA^\dagger$ . all hermitian operators are normal.

**Theorem 1 (Spectral Decomposition Theorem).**

$$M \text{ is a normal operator on } V \iff M \text{ is diagonalisable}$$

diagonalisable w.r.t some basis in  $V$ . The proof is interesting to read, do check it out.

## 1.7 Tensor Products

*Tensor product* is a way of putting vector spaces together to form larger vector spaces. For example, if we have vector spaces  $V$  and  $W$  of dimension  $m$  &  $n$  respectively, the tensor product  $V \otimes W$  represents an  $mn$  dimensional vector space.

Suppose  $|v\rangle \in V$  &  $|w\rangle \in W$  be two vectors, their tensor product  $|v\rangle \otimes |w\rangle$  is also represented as  $|v\rangle|w\rangle$ ,  $|v,w\rangle$  or even  $|vw\rangle$ , and it lies in  $mn$  dimensional space  $V \otimes W$

They have few nice properties:

1. For a scalar  $z$ ,

$$z(|v\rangle \otimes |w\rangle) = (z|v\rangle) \otimes |w\rangle \quad (1.2)$$

2. For  $|v_1\rangle, |v_2\rangle \in V$

$$(|v_1\rangle + |v_2\rangle) \otimes |w\rangle = |v_1\rangle \otimes |w\rangle + |v_2\rangle \otimes |w\rangle \quad (1.3)$$

3. For  $|w_1\rangle, |w_2\rangle \in W$

$$|v\rangle \otimes (|w_1\rangle + |w_2\rangle) = |v\rangle \otimes |w_1\rangle + |v\rangle \otimes |w_2\rangle \quad (1.4)$$

4. If we have operators  $A : V \rightarrow V$  &  $B : W \rightarrow W$ , we can define a new operator  $A \otimes B : V \otimes W \rightarrow V \otimes W$ , we can't write it in matrix form, but we know it's a linear operator and,

$$(A \otimes B)(|v\rangle \otimes |w\rangle) \equiv A|v\rangle \otimes B|w\rangle \quad (1.5)$$

5. We can define inner product in  $V \otimes W$  as

$$(a_i|v_i\rangle \otimes |w_i\rangle, b_j|v_j\rangle \otimes |w_j\rangle) = a_i^*b_j \langle v_i|v_j\rangle \langle w_i|w_j\rangle \quad (1.6)$$

As we have a well defined inner product, this is a hilbert space too, we can use properties like adjoint, unitary, normal, hermitian matrices.

### 1.7.1 Kronecker Product

There is indeed a representation which would let us interpret tensor product in some manner, suppose we have two matrices  $A_{m \times n}, B_{p \times q}$ , their tensor product in this representation is:

$$A \otimes B = \begin{bmatrix} A_{11}B & A_{12}B & \cdots & A_{1n}B \\ A_{21}B & A_{22}B & \cdots & A_{2n}B \\ \vdots & \vdots & \vdots & \vdots \\ A_{m1}B & A_{m2}B & \cdots & A_{mn}B \end{bmatrix}_{mp \times nq} \quad (1.7)$$

for example, tensor product of Pauli matrices  $X$  &  $Y$  is represented as

$$X \otimes Y = \begin{bmatrix} 0 \cdot Y & 1 \cdot Y \\ 1 \cdot Y & 0 \cdot Y \end{bmatrix} = \begin{bmatrix} 0 & 0 & 0 & -i \\ 0 & 0 & i & 0 \\ 0 & -i & 0 & 0 \\ i & 0 & 0 & 0 \end{bmatrix} \quad (1.8)$$

and another example,

$$\begin{bmatrix} 1 \\ 2 \end{bmatrix} \otimes \begin{bmatrix} 2 \\ 3 \end{bmatrix} = \begin{bmatrix} 2 \\ 3 \\ 4 \\ 6 \end{bmatrix} \quad (1.9)$$

Using similar notations,  $|\phi\rangle^{\otimes k}$  means tensor product of  $|\phi\rangle$  with itself  $k$  times.

Just for practice, I'm writing down few tensor products,  $I, X, Y, Z$  are pauli matrices.

$$X \otimes Z = \begin{bmatrix} 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & -1 \\ 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \end{bmatrix} \quad I \otimes X = \begin{bmatrix} X & 0 \\ 0 & X \end{bmatrix} = \begin{bmatrix} 0 & 1 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \end{bmatrix} \quad (1.10)$$

$$X \otimes I = \begin{bmatrix} 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \\ 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \end{bmatrix} \quad (1.11)$$

Hence, tensor product **isn't commutative**. Few more facts,

1.  $(A \otimes B)^* = A^* \otimes B^*$
2.  $(A \otimes B)^T = A^T \otimes B^T$
3.  $(A \otimes B)^\dagger = A^\dagger \otimes B^\dagger$
4. Tensor product preserves (a) Unitarity (b) Hermiticity (c) Positivity (d) Projection

## 1.8 Operator functions

If we have a function  $f : \mathbb{C}^n \rightarrow \mathbb{C}^n$ , we can define a corresponding matrix function, which acts on *normal matrices* as follows

$$A = \sum_a a |a\rangle \langle a| \implies f(A) \equiv \sum_a f(a) |a\rangle \langle a| \quad (1.12)$$

where  $\sum_a a |a\rangle \langle a|$  is a spectral decomposition (exists since  $A$  is normal 1). Now we can calculate square root, logarithm, exponential of a positive, positive-definite, normal operator respectively. for eg.

$$\exp(\theta Z) = \begin{bmatrix} e^\theta & 0 \\ 0 & e^{-\theta} \end{bmatrix} \quad (1.13)$$

I'll do another example, let  $A = \begin{bmatrix} 4 & 3 \\ 3 & 4 \end{bmatrix}$ , let's find square root of it. Eigenvalues of  $A$  are 7, 1. Their corresponding eigenvectors which form an orthonormal basis of  $\mathbb{R}^2$  are  $(|0\rangle + |1\rangle)/\sqrt{2} = |a\rangle$  &  $(|0\rangle - |1\rangle)/\sqrt{2} = |b\rangle$ ,

$$A = 7 |a\rangle \langle a| + |b\rangle \langle b| \quad (1.14)$$

$$\implies \sqrt{A} = \sqrt{7} |a\rangle \langle a| + |b\rangle \langle b| \quad (1.15)$$

$$= \begin{bmatrix} \frac{\sqrt{7}+1}{2} & \frac{\sqrt{7}-1}{2} \\ \frac{\sqrt{7}-1}{2} & \frac{\sqrt{7}+1}{2} \end{bmatrix} \quad (1.16)$$

### 1.8.1 Trace

This is just the sum of diagonal elements of a matrix.

$$\text{tr}(A) = \sum_i A_{ii} \quad (1.17)$$

It's also defined as

$$\text{tr}(A) = \sum_i \langle i | A | i \rangle \quad (1.18)$$

i.e it doesn't depend on which coordinate I represent A on, traces have these nice properties:

1. It's *cyclic*,  $\text{tr}(AB) = \text{tr}(BA)$
2.  $\text{tr}(cA + dB) = c\text{tr}(A) + d\text{tr}(B), c, d \in \mathbb{C}$
3. if  $U$  is unitary,  $\text{tr}(A) = \text{tr}(UAU^\dagger)$
4.  $\text{tr}(A |\psi\rangle \langle \psi|) = \langle \psi | A | \psi \rangle$

### 1.8.2 Hilbert-Schmidt inner product

The vector space  $L_V$  of linear operators on a Hilbert space  $V$  can be converted into a Hilbert space by defining an inner product on  $L_V \times L_V$  as

$$(A, B) \equiv \text{tr}(A^\dagger B) \quad (1.19)$$

this is known as the *Hilbert-Schmidt* or *Trace* inner product.

## 1.9 Commutator & Anti-commutator

The commutator between two operators  $A$  &  $B$  is defined as

$$[A, B] = AB - BA \quad (1.20)$$

if  $[A, B] = 0 \implies AB = BA$ , then we say  $A$  commutes with  $B$ . The *anti-commutator* between  $A$  &  $B$  is defined as

$$A, B = AB + BA \quad (1.21)$$

we say  $A$  anti-commutes with  $B$  if  $A, B = 0$ .

**Theorem 2 (Simultaneous diagonalization theorem).** If  $A$  &  $B$  are Hermitian operators, then  $[A, B] = 0 \implies \exists$  an orthonormal basis such that both  $A$  and  $B$  are simultaneously diagonalizable w.r.t same basis.

### Commutations & Anti-commutation for Pauli matrices

$$[X, Y] = 2iZ \quad [Y, Z] = 2iX \quad [Z, X] = 2iY \quad (1.22)$$

$$(1.23)$$

we can write it using  $\epsilon_{jkl}$ , the antisymmetric tensor, for which  $\epsilon_{jkl} = 0$  except for  $\epsilon_{123} = \epsilon_{231} = \epsilon_{312} = 1$  and  $\epsilon_{321} = \epsilon_{213} = \epsilon_{132} = -1$ :

$$[\sigma_j, \sigma_k] = 2i \sum_{l=1}^3 \epsilon_{jkl} \sigma_l \quad (1.24)$$

Also,  $\{\sigma_i, \sigma_j\} = 0, \forall i \neq j, i, j \in 1, 2, 3$ . whereas  $\{\sigma_0, \sigma_j\} = 2\sigma_j$ .

## 1.10 Polar and Singular value decompositions

**Theorem 3 (Polar decomposition).** Let  $A$  be a linear operator on a vector space  $V$ . Then there exist unitary  $U$  and **positive operators**  $J$  and  $JK$  such that

$$A = UJ = KU \quad (1.25)$$

where the unique positive operators  $J$  &  $K$  are  $J \equiv \sqrt{A^\dagger A}$  and  $K \equiv \sqrt{AA^\dagger}$ , if  $A$  is invertible,  $U$  is unique.

$A = UJ$  is the left polar decomposition of  $A$ , and  $A = KU$  is the right polar decomposition of  $A$

Singular value decomposition is just polar decomposition & spectral decomposition.

**Corollary 3.1 (Singular value decomposition).** Let  $A$  is a square matrix, Then  $\exists$  unitary  $U$  &  $V$ , a diagonal matrix  $D$  with **non-negative entries** such that

$$A = UDV \quad (1.26)$$

diagonal elements of  $D$  are called **singular values** of  $A$ .

*Proof.* By polar decomposition,  $A = SJ$ , where  $S$  is unitary and  $J$  is positive. By spectral decomposition,  $J = TDT^\dagger$  where  $T$  is unitary and  $D$  has non negative entries. setting  $U = ST$  and  $V = T^\dagger$  we get the above thing.  $\square$

# Chapter 2

## Postulates of Quantum Mechanics

### 2.1 State Space

**Postulate 2.1.** Any isolated physical system has a Hilbert Space associated with it known as the **state space**. The system is completely described by its **state vector**, which is a unit vector in state space.

The simplest quantum mechanical system, which we'll be concerned the most about, is a *qubit* which has a 2-dimensional state space. Hence an arbitrary state vector  $|\psi\rangle$  can be represented by

$$|\psi\rangle = \alpha |0\rangle + \beta |1\rangle \quad (2.1)$$

where  $\alpha, \beta \in \mathbb{C}$  and  $|0\rangle, |1\rangle$  form an orthonormal basis of state space. Since state vector is a unit vector,  $\langle\psi|\psi\rangle = 1 \implies |\alpha|^2 + |\beta|^2 = 1$ , this is known as *normalization* of state vector.

For now, *qubit* is an abstract thing. We consider a fixed orthonormal basis  $|0\rangle, |1\rangle$  apriori. These two states can be considered analogous to the bits 0 and 1, except that a state vector is a linear combination, or in other words, *superposition* of these bits. The linear combination  $\sum_i \alpha_i |\psi_i\rangle$  is defined as the superposition of  $|\psi_i\rangle$  with amplitude  $\alpha_i$  for  $|\psi_i\rangle$ .

### 2.2 Evolution

**Postulate 2.2.** The evolution of a **closed** system is described by a **unitary transformation**. The state of the system  $|\psi\rangle$  at time  $t_1$  is related to the state of the system  $|\psi'\rangle$  at time  $t_2$  solely by the unitary transformation  $U$  which only depends on  $t_1, t_2$

$$|\psi'\rangle = U |\psi\rangle \quad (2.2)$$

Quantum mechanics doesn't help us find the *state space* or *unitary operator*, it just assures us that any physical system would behave this way.

Few such operator on a qubit are pauli matrices. X is referred to as *NOT gate* or *bit-flip* since it turns  $|0\rangle$  to  $|1\rangle$  and vice versa. Z is known as *phase-flip* as it leaves  $|0\rangle$  as it is, which inverting the phase (sign) of  $|1\rangle$ .

The second postulate just defines state of the system at discrete intervals  $t_1$  and  $t_2$ , this can be generalised for *continuous* t.

**Corollary 3.2.** *Time evolution of a **closed** quantum system is described by **Schrödinger equation**,*

$$i\hbar \frac{d|\psi\rangle}{dt} = H|\psi\rangle \quad (2.3)$$

where  $\hbar$  is the reduced planck's constant,  $H$  is a fixed **Hermitian operator** known as the **Hamiltonian** of the closed system.

Since Hamiltonian is a Hermitian operator, it has a spectral decomposition,

$$H = \sum_E E |E\rangle \langle E| \quad (2.4)$$

Here the eigenvectors  $|E\rangle$  are conventionally called *energy eigenstates* or *stationary states* and  $E$  is called *energy* of the state  $|E\rangle$ . The lowest energy is known as *ground state energy* and the corresponding state is the *ground state*.

We can see a connection between corollary 2.2 and postulate 3.2, if we know the solution of Schrödinger's equation, which can be verified that it is

$$|\psi(t_2)\rangle = \exp\left[\frac{-iH(t_2 - t_1)}{\hbar}\right] |\psi(t_1)\rangle = U(t_1, t_2) |\psi(t_1)\rangle \quad (2.5)$$

if we set,

$$U(t_1, t_2) \equiv \exp\left[\frac{-iH(t_2 - t_1)}{\hbar}\right] \quad (2.6)$$

it can be shown that if  $K$  is a Hermitian operator, then  $U = \exp(iK)$  is a unitary operator. Thus, there's a one-to-one correspondence between the theorem and it's corollary.

If the system we're considering is not closed, we can approximate it's evolution using a time varying Hamiltonian. There's this example that when we consider a laser focused on an atom, the whole system is properly described by a Hamiltonian, but the behaviour of the atome alone, is seem to be approximately described by another Hamiltonian, this thing contains terms related to the laser intensity.

## 2.3 Quantum measurement

We've seen that in closed systems, the evolution of state vector is according to a unitary transform, but when we try to measure the state of the system, we need to interact with it, which leads to a change other than a unitary transform, which is described as

**Postulate 2.3.** *A **quantum measurement** is a collection  $\{M_m\}$  of measurement operators. These operators act on the state space. The index  $m$  refers to the measurement outcome that we want using  $M_m$ . If the system is in state vector  $|\psi\rangle$  before the measurement had been made, the probability that outcome is  $m$  is,*

$$p(m) = \langle\psi| M_m^\dagger M_m |\psi\rangle \quad (2.7)$$

After the measurement, the state changes to

$$\frac{M_m |\psi\rangle}{\sqrt{\langle\psi| M_m^\dagger M_m |\psi\rangle}} \quad (2.8)$$

All the  $M_i$  satisfy the **completeness relation**

$$\sum_i M_i^\dagger M_i = I \quad (2.9)$$

This ensures the probability that some measurement occurs is 1,

$$\sum_i p(i) = \sum_i \langle \psi | M_i^\dagger M_i | \psi \rangle = 1 \quad (2.10)$$

As an example, suppose we have a closed system whose state vector  $|\psi\rangle$  is currently  $a|0\rangle + b|1\rangle$  and we want to make measurements using  $M_0 = |0\rangle\langle 0|$  and  $M_1 = |1\rangle\langle 1|$ , we can calculate and see that probability that measurements 0 & 1 are made correct are  $|a|^2$  and  $|b|^2$  respectively, and the final states are  $\frac{a}{|a|}|0\rangle$  if we make measurement 0 and  $\frac{b}{|b|}|1\rangle$  if we make measurement 1.

So when there's an isolated physical system, it has a state vector, when you try to measure it, using the collection of measurement operators gives just a label. we can just get the probability that we get a specific probability, using equation 2.7

## 2.4 Distinguishing quantum states

When we have  $n$  systems, with  $n$  state vectors, we can only distinguish them by measuring them only if all the state vectors are orthonormal. This has a good proof.

**Theorem 4.** *No measurement distinguishing non-orthogonal states is possible*

*Proof.* We consider two states  $|\psi_1\rangle, |\psi_2\rangle$  and assume that there exists a measurement. We distinguish these states by looking at the output label we get from the measurement. Taking the function  $f(j)$  where  $j$  is the label. If the state is  $|\psi_1\rangle$  then probability that our prediction is correct  $\sum_{j;f(j)=1} \langle \psi_1 | M_j^\dagger M_j | \psi_1 \rangle = 1$  similarly for 2. To simlify this we define

$$E_i = \sum_{j;f(j)=i} M_j^\dagger M_j \quad (2.11)$$

so from what we've done above

$$\langle \psi_1 | E_1 | \psi_1 \rangle = 1 \langle \psi_2 | E_2 | \psi_2 \rangle = 1 \quad (2.12)$$

since we'll always measure  $|\psi_1\rangle$ 's label when the state is  $|\psi_1\rangle$ ,

$$\langle \psi_1 | E_2 | \psi_2 \rangle = 0 \quad (2.13)$$

$$\implies \sqrt{E_2} | \psi_2 \rangle = 0 \quad (2.14)$$

now let's focus on the orthogonal component of  $|\psi_2\rangle$ , by writing it as  $|\psi_2\rangle = \alpha|\psi_1\rangle + \beta|\phi\rangle$ , where  $|\phi\rangle$  is orthonormal to  $|\psi_1\rangle \implies |\beta| < 1$ , since they're non-orthogonal.

But there's a contradiction,

$$\langle \psi_2 | E_2 | \psi_2 \rangle = \beta^2 \langle \phi | E_2 | \phi \rangle \quad (2.15)$$

$$\leq |\beta|^2 \quad (2.16)$$

$$< 1 \quad (2.17)$$

Thus by contradiction, our assumption is false  $\implies$  we can't distinguish non-orthogonal states reliably.  $\square$

## 2.5 Projective measurements

This is the main thing we'll deal in quantum information and computation.

**Postulate 2.4.** A projective measurement is an observable  $M$ , a Hermitian operator on the state space of system considered, having spectral decomposition

$$M = \sum_m m P_m \quad (2.18)$$

where  $m$  is one of the possible outcomes.  $m$  is an eigenvalue of  $M$ ,  $P_m$  is the projector onto the corresponding eigenspace. When we measure a state  $|\psi\rangle$ , the probability of getting the result  $m$  is

$$p(m) = \langle\psi| P_m |\psi\rangle \quad (2.19)$$

After measuring, the state of the system becomes

$$\frac{P_m |\psi\rangle}{\sqrt{p(m)}} \quad (2.20)$$

We can see projective measurement is a special kind of measurement, which do obey completeness relation, i.e  $\sum_m M_m^\dagger M_m = I$ . But  $M_m$  also are *Orthogonal projectors*, hence *Hermitian* and  $M_m M_{m'} = \delta_{mm'} M_m$

These have few nice properties, we can find mean of measurements easily, which is

$$E(M) = \sum_m mp(m) \quad (2.21)$$

$$= \sum_m m \langle\psi| P_m |\psi\rangle \quad (2.22)$$

$$= \langle\psi| \sum_m m P_m |\psi\rangle \quad (2.23)$$

$$= \langle\psi| M |\psi\rangle \quad (2.24)$$

this is a nice thing, we can also find the standard deviation of the outcomes we get.

$$\Delta(M) = \sqrt{\langle(M - \langle M \rangle)^2\rangle} \quad (2.25)$$

$$= \sqrt{\langle M^2 \rangle - \langle M \rangle^2} \quad (2.26)$$

As you might have observed, we defined a new term *observable*, some still like to look at the projection measurement as a collection of operators, using the observable implicitly. Another phrase used is “measure in basis  $|m\rangle$ ”, which means the projectors  $P_m = |m\rangle \langle m|$  are used for projective measurement.

If  $\vec{v}$  be a three dimensional unit vector, then the measurement of the observable

$$\vec{v} \cdot \vec{\sigma} = v_1 \sigma_1 + v_2 \sigma_2 + v_3 \sigma_3 \quad (2.27)$$

is historically known as the “*measurement of spin along  $\vec{v}$  axis*”

*Remark.* Eigenvalues of  $\vec{v} \cdot \vec{\sigma}$  are  $\pm 1$  and the corresponding projection operators are  $P_\pm = \frac{I \pm \vec{v} \cdot \vec{\sigma}}{2}$

### 2.5.1 Heisenberg Uncertainty Principle

Suppose  $A, B$  be two Hermitian operators and  $|\psi\rangle$  is a quantum state. Let  $\langle\psi|AB|\psi\rangle = x + iy$  where  $x, y \in \mathbb{R}$ . Then  $\langle\psi|[AB]|\psi\rangle = 2iy$  and  $\langle\psi|\{AB\}|\psi\rangle = 2x$ , therefore

$$(\langle\psi|[A, B]|\psi\rangle)^2 + (\langle\psi|\{A, B\}|\psi\rangle)^2 = 4|\langle\psi|AB|\psi\rangle|^2 \quad (2.28)$$

Using Cauchy-Schwartz inequality, then substituting  $A = C - \langle C \rangle$  and  $B = D - \langle D \rangle$

$$\langle\psi|A^2|\psi\rangle\langle\psi|B^2|\psi\rangle \geq (\langle\psi|AB|\psi\rangle)^2 \quad (2.29)$$

$$\geq \frac{(\langle\psi|[A, B]|\psi\rangle)^2}{4} \quad (2.30)$$

$$\Rightarrow \Delta(C)\Delta(D) \geq \frac{\langle\psi|[C, D]|\psi\rangle}{2} \quad (2.31)$$

Hence Proved.

Note that the common misconception of Heisenberg uncertainty principle is that we “measure a state with some error” which is  $\Delta(C), \Delta(D)$ , but the actual thing is we have a large amount of  $|\psi\rangle$ ’s and we measure each of them separately to with observables  $C, D$  to get a set of values of  $C, D$  whose standard deviation is  $\Delta(C), \Delta(D)$ . This is because once we make a measurement, the quantum state changes because we’ve altered it.

## 2.6 POVM measurements

It stands for *Positive Operator-Valued Measure*. POVM is a formalism, and is used when we are concerned with the outcome more than the state of the output we’re measuring. For eg. when we’ve done an experiment and just want to measure the final state.

If we’re measuring a state  $|\psi\rangle$  using measurement operators  $M_m$ , the probability that the output is  $m$  is  $\langle\psi|M_m^\dagger M_m|\psi\rangle$ , if we define

$$E_m = M_m^\dagger M_m \quad (2.32)$$

$E_m$  satisfy  $\sum_m E_m = I$ ,  $E_m$  is enough to find the probability of output  $m$  which is  $\langle\psi|E_m|\psi\rangle$ . The set  $\{E_m\}$  is called the *POVM set*.

In projective measurements POVM set is same as the set of measurement operators, since  $E_m = P_m^\dagger P_m = P_m$ . In fact any measurement where the measurement operators and the POVM elements coincide is a projective measurement.

If we want the post-measurement state, we can conveniently *define* a POVM to be any set of operators  $\{E_m\}$  such that (a) each  $E_m$  is positive (b) it satisfies *completeness relation*, i.e  $\sum_m E_m = I$ . In this way we can choose  $M_m = U_m \sqrt{E_m}$  where  $U_m$  is a unitary operator.

## 2.7 Phase

This is kind of interesting and deja vu giving.

### 2.7.1 Global Phase

If  $|\psi\rangle$  is a state,  $\theta \in \mathbb{R}$ , we say that  $e^{i\theta}|\psi\rangle$  is equal to  $|\psi\rangle$  upto the *global phase factor*  $e^{i\theta}$ . These two act quite similar when you try to measure them, because the output of m would occur with a probability of  $\langle\psi|M_m^\dagger M_m|\psi\rangle$  and  $\langle\psi|e^{-i\theta}M_m^\dagger M_m e^{i\theta}|\psi\rangle = \langle\psi|M_m^\dagger M_m|\psi\rangle$ , which is the same. Thus, they’re physically equivalent.

### 2.7.2 Relative Phase

If two amplitudes,  $a$  and  $b$ , differ by a *relative phase* if there is a real  $\theta$  such that  $a = be^{i\theta}$ . Two states differ by a relative phase in some basis if each of the amplitudes in that basis is related by such a phase factor. Unlike global phase factors, relative phase factors differ from amplitude to amplitude, thus, there are physically observable differences in measurement statistics, and two states differing by a relative phase aren't physically equivalent.

## 2.8 Composite systems

It's useful in describing state of a composite system made up of multiple systems

**Postulate 2.5.** *The state space of a composite physical system is the tensor product of state spaces of component physical systems. If  $n$  physical systems are prepared in state  $|\psi_1\rangle, \dots, |\psi_n\rangle$  The composite system has the state  $|\psi_1\rangle \otimes |\psi_2\rangle \otimes \dots \otimes |\psi_n\rangle$*

Using this, we can prove a fact.

*Remark.* To implement a general measurement, it is sufficient to make projective measurements together with unitary dynamics

I'm not giving the entire proof. but assume that the state space we're considering is  $Q$  and we want to measure using operators  $M_m$ . We can consider a new *ancilla* system which has state space  $M$ .  $M$  has orthonormal basis  $|m\rangle$  which are in one-to-one correspondence with the first output. Let's define an operator  $U : QM \rightarrow QM$  as

$$U|\psi\rangle|0\rangle = \sum_m M_m|\psi\rangle|m\rangle \quad (2.33)$$

where  $|0\rangle$  is some state in  $Q$ . It can be shown that  $U$  preserves dot product. i.e  $\langle 0 | \langle \phi | U^\dagger U |\psi\rangle |0\rangle = \langle \phi | \psi \rangle$  which implies that  $U$  extends over  $Q \otimes M$  which can be proved too.

Now, if we consider two systems with projectors  $P_m = I_Q \otimes |m\rangle\langle m|$  and measure using unitary dynamics and projective transformation, the probability of output  $m$  is

$$p(m) = \langle \psi | \langle 0 | U^\dagger P_m U | 0 \rangle | \psi \rangle \quad (2.34)$$

$$= \sum_{m',m''} \langle \psi | \langle m' | M_{m'}^\dagger (I_Q \otimes |m\rangle\langle m|) M_{m''} | m'' \rangle | \psi \rangle \quad (2.35)$$

$$= \langle \psi | M_m^\dagger M_m | \psi \rangle \quad (2.36)$$

which is just as measured directly from general measurement. also the state of the composite system after the measurement is

$$\frac{P_m U |\psi\rangle |0\rangle}{\sqrt{\langle \psi | U^\dagger P_m U | \psi \rangle}} = \frac{M_m |\psi\rangle |m\rangle}{\sqrt{\langle \psi | M_m^\dagger M_m | \psi \rangle}} \quad (2.37)$$

Since the state of the system  $M$  after measurement is  $|m\rangle$ , state of system  $Q$  is

$$\frac{M_m |\psi\rangle}{\sqrt{\langle \psi | M_m^\dagger M_m | \psi \rangle}} \quad (2.38)$$

also as measured directly from general measurement.

## Quantum Entanglement

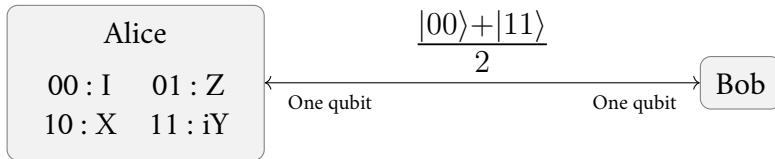
This is something you would've heard of before. Consider the two qubit state,

$$|\psi\rangle = \frac{|00\rangle + |11\rangle}{\sqrt{2}} \quad (2.39)$$

we can never write it as a tensor product of two other qubit states, i.e  $|\psi\rangle = |a\rangle|b\rangle$  can never be true. It's easily checkable. A state of a composite system like this, which can't be written as tensor product of its component systems is an *entangled* state. Entangled states play a crucial role in quantum computation and information.

## 2.9 Application: Superdense Coding

This nicely uses the idea of entanglement and “applying” quantum gates. Suppose we have ‘Alice’ and ‘Bob’ again, who’re away from each other. Alice wants to share two classical bits of information using one qubit to Bob, can she do it?



Yes, using superdense coding. Here, Alice and Bob share a pair of externally setup qubits in entangled state

$$|\psi\rangle = \frac{|00\rangle + |11\rangle}{\sqrt{2}} \quad (2.40)$$

Alice has one qubit, which she can alter and Bob has another. Any change Alice makes to her qubit changes the composite system, which can be measured by Bob. So she applies quantum gates to change the qubits, as given in the diagram, which result in the following changes.

$$00 : |\psi\rangle \xrightarrow{I} \frac{|00\rangle + |11\rangle}{\sqrt{2}} \quad (2.41)$$

$$01 : |\psi\rangle \xrightarrow{Z} \frac{|00\rangle - |11\rangle}{\sqrt{2}} \quad (2.42)$$

$$10 : |\psi\rangle \xrightarrow{X} \frac{|10\rangle + |01\rangle}{\sqrt{2}} \quad (2.43)$$

$$11 : |\psi\rangle \xrightarrow{iY} \frac{-|10\rangle + |01\rangle}{\sqrt{2}} \quad (2.44)$$

We can see that all the outputs are orthonormal, hence Bob can choose appropriate measurement operators and clearly distinguish each state, hence getting the information Alice had sent him. These are known as the *Bell states*, *Bell basis* or *EPR pairs*.

Thus, a single qubit can be used to transfer information which would've taken two classical bits.

## 2.10 The density operator

Till now we've been using the language of state vectors. The density operator is a different way of representing states, it's mathematically equivalent to state vectors. It's much more convenient to use in some scenarios. It's really useful while describing *individual subsystems* of a composite system.

### 2.10.1 Ensembles of quantum states

It's useful when we don't exactly know the state of a system. If the possible states are  $|\psi_i\rangle$  with probabilities  $p_i$ ,  $\{p_i, |\psi_i\rangle\}$  is an *ensemble of pure states*. The density operator, also called as the *density matrix* is defined as

$$\rho \equiv \sum_i p_i |\psi_i\rangle \langle \psi_i| \quad (2.45)$$

We can now workout rewriting the postulates of quantum mechanics.

Suppose, the evolution of closed quantum system is described by a unitary operator  $U$ . The system is initially in the state  $|\psi_i\rangle$  with probability  $p_i$ , after the evolution, system will be in the state  $U|\psi_i\rangle$  with probability  $p_i$ , then the density operator evolves into

$$\rho = \sum_i p_i |\psi_i\rangle \langle \psi_i| \xrightarrow{U} \sum_i p_i U |\psi_i\rangle \langle \psi_i| U^\dagger = U \rho U^\dagger. \quad (2.46)$$

We can also predict the probability of a measurement outcome  $m$ , when our system is defined by the density operator  $\rho$ , if the initial state was  $|\psi_i\rangle$ , then the probability of outcome  $m$  is

$$p(m|i) = \langle \psi_i | M_m^\dagger M_m | \psi_i \rangle = \text{tr}(M_m^\dagger M_m |\psi_i\rangle \langle \psi_i|) \quad (2.47)$$

This is the probability that the outcome would be  $m$ , if the state was  $|\psi_i\rangle$  which has a probability of  $p_i$ , hence the overall probability of outcome being  $m$  is

$$p(m) = \sum_i p(m|i)p_i \quad (2.48)$$

$$= \sum_i p_i \text{tr}(M_m^\dagger M_m |\psi_i\rangle \langle \psi_i|) \quad (2.49)$$

$$= \text{tr}(M_m^\dagger M_m \rho) \quad (2.50)$$

Also we can find out the state of the system if the outcome was  $m$ , if the state was initially  $|\psi_i\rangle$  which has probability  $p_i$  the final state would be

$$|\psi_i^m\rangle = \frac{M_m |\psi_i\rangle}{\sqrt{\langle \psi_i | M_m^\dagger M_m | \psi_i \rangle}} \quad (2.51)$$

Hence, the final density operator would be

$$\rho_m = \sum_i p(i|m) |\psi_i^m\rangle \langle \psi_i^m| = \sum_i p(i|m) \frac{M_m |\psi_i\rangle \langle \psi_i| M_m^\dagger}{\langle \psi_i | M_m^\dagger M_m | \psi_i \rangle} \quad (2.52)$$

By Bayes' theorem,  $p(i|m) = p(m|i)p_i/p(m)$ , substituting this,

$$\rho_m = \sum_i p_i \frac{M_m |\psi_i\rangle \langle \psi_i| M_m^\dagger}{\text{tr}(M_m^\dagger M_m \rho)} \quad (2.53)$$

$$= \frac{M_m \rho M_m^\dagger}{\text{tr}(M_m^\dagger M_m \rho)} \quad (2.54)$$

A system whose state  $|\psi\rangle$  is exactly known is said to be in *pure state*. The density operator is then simply  $|\psi\rangle \langle \psi|$ . Otherwise, the system is said to be in a *mixed state*. It's said to be in a *mixture* of different states in the ensemble of  $\rho$ . Suppose a quantum state is prepared from  $\rho_i$  with probability  $p_i$ , the system can then be described by the density operator  $\sum_i p_i \rho_i$ , which can be proved.

This is useful mainly in the analysis of quantum noise, when we need to introduce ignorance into our knowledge of the quantum state. For example if we made a measurement and somehow lost/forgot our measurement, then the final state of the system would be *mixed*, given by

$$\rho = \sum_m p(m) \rho_m \quad (2.55)$$

$$= \sum_m \text{tr}(M_m^\dagger M_m \rho) \frac{M_m \rho M_m^\dagger}{\text{tr}(M_m^\dagger M_m \rho)} \quad (2.56)$$

$$= \sum_m M_m \rho M_m^\dagger \quad (2.57)$$

## 2.11 General properties of density operator

We can actually characterize what operators are density operators using

**Theorem 5 (Characterization of density operators).** *An operator  $\rho$  is a density operator associated with an ensemble  $\{p_i, |\psi_i\rangle\}$  if and only if*

1. **Trace condition**  $\text{tr}(\rho) = 1$
2. **Positivity condition**  $\rho$  is a positive operator.

*Proof.* Forward is easy and upto the reader. For the backward proof, a possible ensemble is  $\{\lambda_i, |i\rangle\}$ , where  $\lambda_i, |i\rangle$  are eigenvalues and corresponding eigenvectors of  $\rho$ , it's upto the reader to check other conditions.  $\square$

With this in our hands, we can redefine our postulates with this density operator picture

**Postulate 1.** *Every physical system is associated with a Hilbert space known as its state space. The system is completely described by its density operator,  $\rho$  which is a positive operator with trace equal to one acting on the state space of the system. If the system is in state  $\rho_i$  with probability  $p_i$ , the density operator of the system is  $\sum_i p_i \rho_i$ .*

**Postulate 2.** *The evolution of a closed system is described by a unitary transformation, i.e state of the system  $\rho$  at time  $t_1$  is related to the state  $\rho'$  at time  $t_2$  by a unitary operator which only depends on  $t_1$  and  $t_2$ .*

$$\rho' = U \rho U^\dagger \quad (2.58)$$

**Postulate 3.** A quantum measurement is described by a set  $\{M_m\}$  of measurement operators. These act on the state space of the system. The index  $m$  refers to the possible outcome of the measurement. If the state of the system is  $\rho$  just before the measurement, the probability of outcome  $m$  is

$$p(m) = \text{tr}(M_m \rho M_m^\dagger) \quad (2.59)$$

The state of the system changes after the measurement is made, to  $\rho'$ , given by

$$\rho' = \frac{M_m \rho M_m^\dagger}{\text{tr}(M_m \rho M_m^\dagger)} \quad (2.60)$$

The measurement operators satisfy the relation

$$\sum_m M_m M_m^\dagger = I. \quad (2.61)$$

**Postulate 4.** State space of a composite system is the tensor product of its component systems. If the components are prepared in state  $\rho_1, \rho_2, \dots, \rho_n$ , state of the composite system is given by  $\rho_1 \otimes \rho_2 \otimes \dots \otimes \rho_n$ .

*Remark.* If  $\rho$  is a density operator, then  $\text{tr}(\rho^2) \leq 1$ . Equality holds when  $\rho$  is a pure state.

By now, you would've noticed that two *different* states can have the same density operator. To get to know what kind of connection there is between these states, we define  $|\tilde{\psi}_i\rangle$  associated with the state  $|\psi_i\rangle$ , the latter is said to *generate* the density operator if  $\rho = \sum_i |\tilde{\psi}_i\rangle \langle \tilde{\psi}_i|$ . It's observable that  $|\tilde{\psi}_i\rangle = \sqrt{p_i} |\psi_i\rangle$ . Now we proceed for the theorem.

**Theorem 6** (Unitary freedom in the ensemble of the density matrices.). Two sets  $|\tilde{\psi}_i\rangle$  and  $|\tilde{\phi}_j\rangle$  generate the same density operator if

$$|\tilde{\psi}_i\rangle = \sum_j u_{ij} |\tilde{\phi}_j\rangle \quad (2.62)$$

where  $u_{ij}$  is a unitary matrix. The sets which is smaller is padded with 0.

In other words, if they two systems have possible states  $|\psi_i\rangle$  and  $|\phi_j\rangle$  with probabilities  $p_i$  and  $q_j$  respectively, they both have the same density operator if and only if

$$\sqrt{p_i} |\psi_i\rangle = \sum_j u_{ij} \sqrt{q_j} |\phi_j\rangle \quad (2.63)$$

where  $u_{ij}$  is a unitary matrix, and the smaller set is padded with 0.

## 2.12 Reduced Density Operator

This is most useful in describing *subsystems* of a composite quantum system. Suppose we have physical system  $A$  and  $B$ , whose state is described by the density operator  $\rho^{AB}$ . The reduced density operator for system  $A$  is defined by

$$\rho^A \equiv \text{tr}_B(\rho^{AB}), \quad (2.64)$$

where  $\text{tr}_B$  is a map of operators known as the *partial trace* over system  $B$ . It's defined by

$$\text{tr}_B(|a_1\rangle \langle a_2| \otimes |b_1\rangle \langle b_2|) \equiv |a_1\rangle \langle a_2| \text{tr}(|b_1\rangle \langle b_2|), \quad (2.65)$$

This partial trace operation is defined only on a special subclass of operators on  $AB$ ; the specification is completed by requiring in addition to Equation (2.65) that the partial trace be linear in its input.

Let's do an example to understand this better, suppose  $\rho^{AB} = \rho \otimes \sigma$ , where  $\rho$  is the density operator for system A, and  $\sigma$  for system B. Then

$$\rho^A = \text{tr}_B(\rho \otimes \sigma) = \rho \text{tr}(\sigma) = \rho \quad (2.66)$$

Which is intuitively correct. For an unintuitive example, let's take an entangled state which would further reveal many things. Let's take the bell state  $AB \frac{|00\rangle + |11\rangle}{\sqrt{2}}$ , which is an entangled state and has the density operator

$$\rho^{AB} = \left( \frac{|00\rangle + |11\rangle}{\sqrt{2}} \right) \left( \frac{\langle 00| + \langle 11|}{\sqrt{2}} \right) \quad (2.67)$$

$$= \frac{|00\rangle \langle 00| + |11\rangle \langle 00| + |00\rangle \langle 11| + |11\rangle \langle 11|}{2} \quad (2.68)$$

Thus we can get  $\rho^B$  from this

$$\rho^B = \text{tr}_A(\rho^{AB}) \quad (2.69)$$

$$= \frac{\text{tr}(|0\rangle \langle 0|)|0\rangle \langle 0| + \text{tr}(|1\rangle \langle 0|)|1\rangle \langle 0|)}{2} \quad (2.70)$$

$$+ \frac{\text{tr}(|0\rangle \langle 1|)|0\rangle \langle 1| + \text{tr}(|1\rangle \langle 1|)|1\rangle \langle 1|)}{2} \quad (2.71)$$

$$= \frac{|0\rangle \langle 0| + |1\rangle \langle 1|}{2} \quad (2.72)$$

$$= \frac{I}{2}. \quad (2.73)$$

This is actually a *mixed state*, since  $\text{tr}((I/2)^2) < 1$ . It's wonderful how we can't find state vector of an entangled qubit, but we can find out the density operator of an entangled qubit. Also the composite system is in pure state, but the entangled qubit is in a mixed state, we don't know exactly what it is. This is a really strange property, we can know the exact state of a joint system, but still we can't figure out each subsystem's state exactly. A hallmark of quantum entanglement.

A physical justification for making such an identification is that the reduced density operator provides the correct measurement statistics for the measurements made on system A.

## 2.13 Schmidt decomposition and purifications

### 2.13.1 Schmidt decomposition

This is a tool like density operator which are useful for the study of composite quantum systems.

**Theorem 7 (Schmidt decomposition).** Suppose  $|\psi\rangle$  is a pure state of a composite system,  $AB$ . Then there exist orthonormal state  $|i_A\rangle$  for system A, and orthonormal states  $|i_B\rangle$  of system B such that

$$|\psi\rangle = \sum_i \lambda_i |i_A\rangle |i_B\rangle, \quad (2.74)$$

where  $\lambda_i$  are **non-negative real numbers** satisfying  $\sum_i \lambda_i^2 = 1$  known as the **Schmidt coefficients**.

Do note that it's not the linear combination of all possible tensor product of bases of  $A$  and  $B$ . It's just  $i$ , i.e linear. Only the tensor product of corresponding bases. This result is very useful. Let  $|\psi\rangle$  be a pure state of a composite system,  $AB$ . Then by Schmidt decomposition,  $\rho^A = \sum_i \lambda_i^2 |i_A\rangle \langle i_A|$  and  $\rho^B = \sum_i \lambda_i^2 |i_B\rangle \langle i_B|$ . Thus eigenvalues of  $\rho^A$  and  $\rho^B$  are identical, namely  $\lambda_i^2$  for both density operators.

*Proof.* We'll consider when  $A$  and  $B$  have same dimension, it's also true when they're not of same dimension. Let  $|j\rangle$  and  $|k\rangle$  be fixed o.n.b for  $A$  and  $B$  respectively, then  $|\psi\rangle$  can be written as

$$|\psi\rangle = \sum_{jk} a_{jk} |j\rangle |k\rangle, \quad (2.75)$$

where  $a_{jk}$  form a matrix  $a$ . By singular value decomposition,  $a = u d v$ , where  $d$  is a diagonal matrix with non-negative elements,  $u$  and  $v$  are unitary matrices. Thus

$$|\psi\rangle = \sum_{ijk} u_{ji} d_{ii} v_{ik} |j\rangle |k\rangle, \quad (2.76)$$

If we define,  $|i_A\rangle \equiv \sum_j u_{ji} |j\rangle$ ,  $|i_B\rangle \equiv \sum_k v_{ik} |k\rangle$ , and  $\lambda_i \equiv d_{ii}$ , it gives

$$|\psi\rangle = \sum_i \lambda_i |i_A\rangle |i_B\rangle \quad (2.77)$$

$|i_A\rangle$  and  $|i_B\rangle$  form an orthonormal set, from the unitarity of  $u$  and  $v$  and the orthonormality of  $|j\rangle$  and  $|k\rangle$ .  $\square$

*Remark.* If  $|ABC\rangle$  is a three component quantum system. There are quantum states  $|\psi\rangle$  of the system which can't be written as

$$|\psi\rangle = \sum_i \lambda_i |i_A\rangle |i_B\rangle |i_C\rangle, \quad (2.78)$$

where  $\lambda_i \in \mathbb{R}$  and  $|i_A\rangle, |i_B\rangle, |i_C\rangle$  are the orthonormal bases of the respective systems.

The bases  $|i_A\rangle$  and  $|i_B\rangle$  are known as the *Schmidt bases* of systems  $A$  and  $B$ . The number of non-zero values  $\lambda_i$  is known as the *Schmidt number* of state  $|\psi\rangle$ . It in some sense relates to the "amount" of entanglement in  $|\psi\rangle$ . Schmidt number is preserved under unitary transformation, i.e Schmidt number of  $U|\psi\rangle$  = Schmidt number of  $|\psi\rangle$ , since  $U|\psi\rangle = \sum_i \lambda_i (U|i_A\rangle) |i_B\rangle$ .

### 2.13.2 Purification

Suppose we're given a state  $\rho^A$  of a quantum system  $A$ , which might be *mixed* or *pure*. We can surely introduce a new system  $R$  and define a *pure state*  $|AR\rangle$  for  $AR$  such that  $A = \text{tr}_R(|AR\rangle \langle AR|)$ . Here  $R$  is known as the *reference system* has no physical significance, we just introduced it to *purify* our mixed state, i.e associate our mixed state with a pure state. This thing is known as *purification* of  $\rho^A$ .

To prove this, suppose  $\rho^A$  has the orthonormal decomposition  $\sum_i p_i |i^A\rangle \langle i^A|$ . We now introduce a system  $R$  which has the same state space as  $A$ , and has an orthonormal basis  $|i^R\rangle$ , to define a pure state of the combined system  $AR$ ,

$$|AR\rangle \equiv \sum_i \sqrt{p_i} |i_A\rangle |i_R\rangle \quad (2.79)$$

Now the reduced density operator of system  $A$ , corresponding to the state  $|AR\rangle$  is

$$\text{tr}_R(|AR\rangle \langle AR|) = \sum_{ij} \sqrt{p_i p_j} |i^A\rangle \langle j_A| \text{tr}(|i^R\rangle \langle j^R|) \quad (2.80)$$

$$= \sum_{ij} \sqrt{p_i p_j} |i^A\rangle \langle j_A| \delta_{ij} \quad (2.81)$$

$$= \sum_i p_i |i^A\rangle \langle i^A| \quad (2.82)$$

$$= \rho^A \quad (2.83)$$

Thus,  $|AR\rangle$  is the purification of  $\rho^A$ .

## **Part II**

### **Week 2**

# Chapter 3

## Overview of Quantum Computing

### 3.1 Quantum Bits

A *bit* is a fundamental concept in classical computation and classical information. An analogous thing in QIC is a *quantum bit* or *qubit*. We'll treat it as a mathematical object but it also has correspondence with real world objects. Like a classical bit has 0, 1 as its states, a qubit's analogous states are  $|0\rangle$  and  $|1\rangle$ . In general a qubit can have the general state which is a *linear combination* or *superposition* of  $|0\rangle$  and  $|1\rangle$  as

$$|\psi\rangle = \alpha|0\rangle + \beta|1\rangle, \alpha, \beta \in \mathbb{C} \quad (3.1)$$

$|0\rangle, |1\rangle$  are known as *computational basis states*. Unlike a classical bit, we can't determine state of a qubit. When we try observing/measuring it, we get  $|0\rangle$  with probability  $|\alpha|^2$  and  $|1\rangle$  with probability  $|\beta|^2$ . Thus  $|\alpha|^2 + |\beta|^2 = 1$ . Hence we can think of the state as unit vector in a complex vector space with the computational bases. This unobservability of a qubit makes up the heart of QIC.

A state  $|\psi\rangle$  of a qubit can be represented as,

$$|\psi\rangle = e^{i\gamma} \left( \cos \frac{\theta}{2} |0\rangle + e^{i\varphi} \sin \frac{\theta}{2} |1\rangle \right) \quad (3.2)$$

which equal to this by a global factor  $e^{i\gamma}$ ,

$$|\psi\rangle = \cos \frac{\theta}{2} |0\rangle + e^{i\varphi} \sin \frac{\theta}{2} |1\rangle \quad (3.3)$$

$\theta, \varphi$  define a point on a sphere known as the *Bloch sphere*; this thing describes many operations of qubits properly. But this doesn't generalize when it comes to multiple qubits.

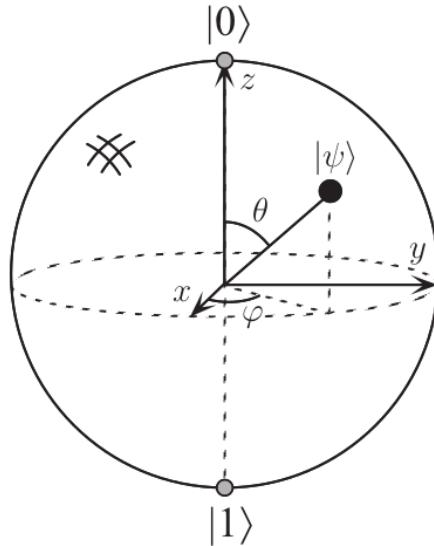


Figure 3.1: Bloch sphere representation of a qubit.

### 3.1.1 Multiple qubits

State vector  $|\psi\rangle$  of a system of two qubits can be defined as,

$$|\psi\rangle = \alpha_{00} |00\rangle + \alpha_{01} |01\rangle + \alpha_{10} |10\rangle + \alpha_{11} |11\rangle \quad (3.4)$$

It can be seen that it has four *computational bases*,  $|00\rangle, |01\rangle, |10\rangle, |11\rangle$ . Even here when we measure the qubit, we get  $x = (00, 01, 10, 11)$  with probability  $|\alpha_x|^2$  and the qubit collapses to the state  $|x\rangle$ . By *normalization*,  $\sum_{x \in \{0,1\}^2} |\alpha_x|^2 = 1$ , where  $\{0, 1\}^2$  is set of all strings with each character 0 or 1. For this set of two qubits, if we try to measure a subset of this, i.e. a qubit, say first qubit, the probability of getting the output 0 is  $|\alpha_{00}|^2 + |\alpha_{01}|^2$  and after measuring, the system collapses to

$$|\psi'\rangle = \frac{\alpha_{00} |00\rangle + \alpha_{01} |01\rangle}{\sqrt{|\alpha_{00}|^2 + |\alpha_{01}|^2}} \quad (3.5)$$

This is a similar collapse *re-normalized* with a normalization constant.

Another interesting thing can be seen if we consider the Bell state,

$$|\psi\rangle = \frac{|00\rangle + |11\rangle}{2} \quad (3.6)$$

If we try to measure  $|0\rangle$ , we get the output 0 with probability  $\frac{1}{2}$  and the bell state collapses to  $|00\rangle$ , and the output 1 with probability  $\frac{1}{2}$  with the bell state collapsing to  $|11\rangle$ . We can see that if we get 0 we perfectly know about second qubit, i.e measurements are *correlated*. Also, a set of  $n$  systems can store  $2^n$  amount of information, even though when we measure it, it collapses.

## 3.2 Quantum Computation

Similar to an electrical circuit, there's a *quantum circuit* which performs operations on qubits using *quantum gates*.

### 3.2.1 Single qubit gates

We can think of a NOT gate similar to its classical analogue, which converts  $|0\rangle$  to  $|1\rangle$  and vice-versa. It's also linear, i.e NOT gate when acted on the state

$$\alpha|0\rangle + \beta|1\rangle \quad (3.7)$$

it interchanges  $|0\rangle$  and  $|1\rangle$  to give

$$\alpha|1\rangle + \beta|0\rangle \quad (3.8)$$

This fact is very important, which can be realised later. NOT gate can be represented by the matrix

$$X \equiv \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix} \quad (3.9)$$

For a matrix  $U$  to be a *quantum gate*, it has to be unitary and this is the *only* constraint.

Few interesting gates are  $Z$ , which leaves  $|0\rangle$  unchanged while flips the sign of  $|1\rangle$ ,

$$Z \equiv \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix} \quad (3.10)$$

The output is equal upto a relative phase factor, the other is the *Hadamard gate*,  $H$  given by

$$H \equiv \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 1 \\ 1 & -1 \end{bmatrix} \quad (3.11)$$

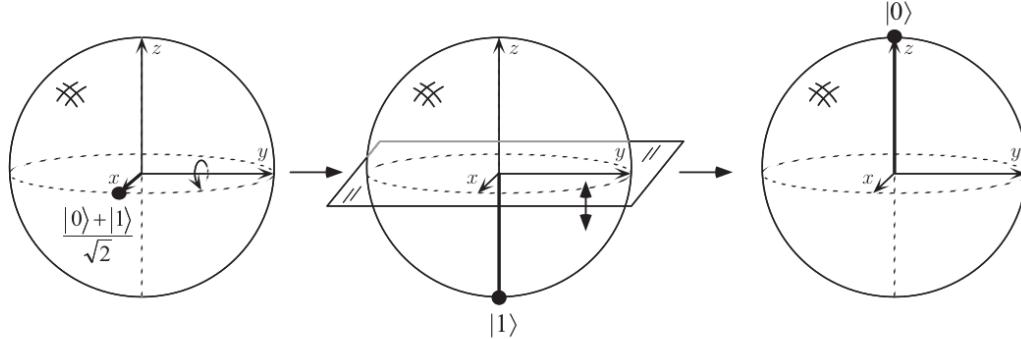


Figure 3.2: Hadmard gate acting on  $\frac{|0\rangle+|1\rangle}{\sqrt{2}}$  to give  $|0\rangle$

This gets takes  $|0\rangle$  to  $\frac{|0\rangle+|1\rangle}{\sqrt{2}}$  and  $|1\rangle$  to  $\frac{|0\rangle-|1\rangle}{\sqrt{2}}$ . In the Bloch sphere point of view, this rotates the sphere first w.r.t  $\hat{y}$  axis by  $90^\circ$  and then rotates w.r.t  $\hat{x}$  axis by  $180^\circ$ .

These qubit gates are summarised as

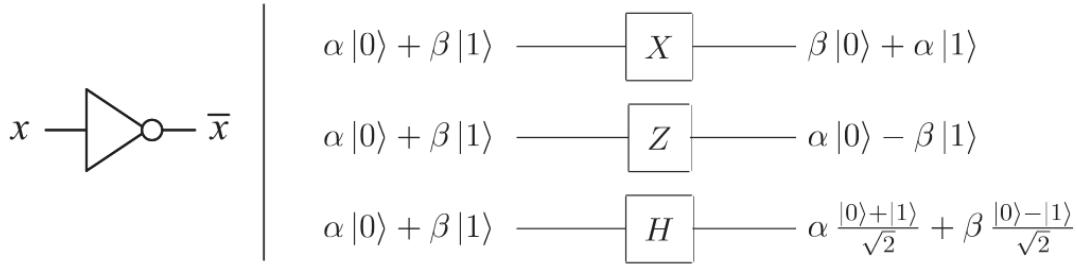


Figure 3.3: Single bit (left) and Qubit (right) logic gates.

We can see that any unitary matrix corresponds to a single qubit gate. Since, there are infinitely many unitary matrices, there are infinitely many single qubit gates. But we can represent this set only with a few single qubit gates. It can be shown that any unitary matrix  $U$  can be written as

$$U = e^{i\alpha} \begin{bmatrix} e^{-i\beta/2} & 0 \\ 0 & e^{i\beta/2} \end{bmatrix} \begin{bmatrix} \cos \frac{\gamma}{2} & \sin \frac{\gamma}{2} \\ \sin \frac{\gamma}{2} & -\cos \frac{\gamma}{2} \end{bmatrix} \begin{bmatrix} e^{-i\delta/2} & 0 \\ 0 & e^{i\delta/2} \end{bmatrix} \quad (3.12)$$

This can be seen as a rotation about  $\hat{z}$  axis then a rotation about  $\hat{y}$  axis followed by a rotation about  $\hat{z}$  axis again and global phase shift. Any single qubit gate can be decomposed into a form like this. Also, any computation on a *finite* number of qubits can be generated by a *finite* set of gates known as the *universal* for quantum computation.

### 3.2.2 Multiple qubit gates

For classical bits, there are AND, OR, NAND, NOR, XOR gates. NAND gate is considered as the *universal gate* as any function of bits can be constructed using NAND gates. XOR gate can't do this. There's a special multi-qubit quantum gate, known as *controlled-NOT* gate or CNOT gate, which has a *control qubit* and a *target qubit*. If the control qubit is 0, nothing happens to target qubit, whereas if it's 1, the target qubit is flipped. i.e,

$$|00\rangle \rightarrow |00\rangle ; |01\rangle \rightarrow |01\rangle ; |10\rangle \rightarrow |11\rangle ; |11\rangle \rightarrow |10\rangle \quad (3.13)$$

The CNOT gate is similar to the XOR gate, because it's just doing  $|A, B\rangle \rightarrow |A, B \oplus A\rangle$ , where  $B \oplus A$  is just the XOR gate, is the sum modulo two. There are no similar generalisations for classical XOR, NAND because they're *irreversible* or *non-invertible*. Information is lost upon the action of these gates which is irretrievable. CNOT and single qubit gates are prototypes for all other gates because of the following *universality* result: *Any multiple qubit gate may be composed from CNOT and single qubit gates.*

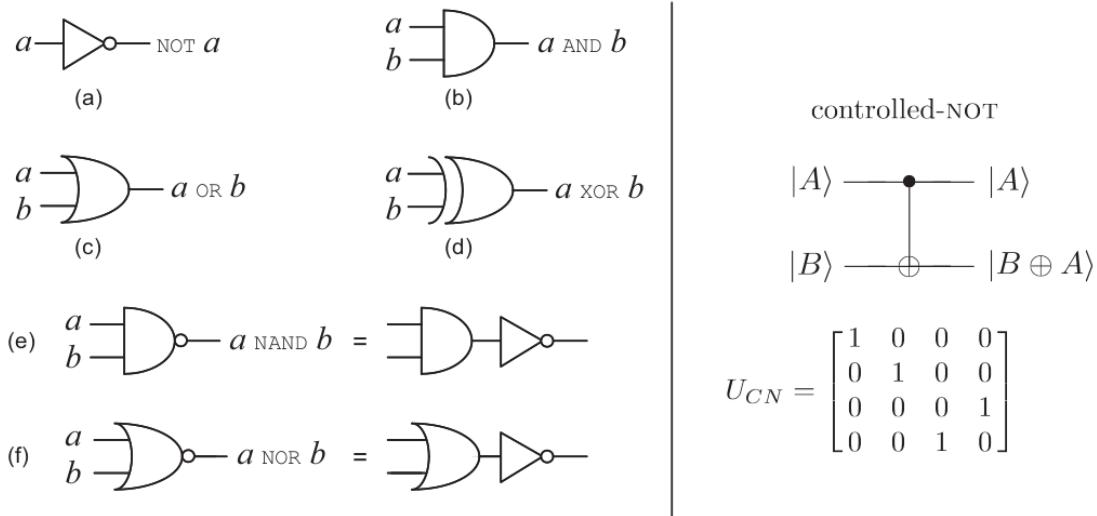


Figure 3.4: Figure on left shows common bit gates; The right figure shows the CNOT gate along with its matrix representation, in the order of  $|00\rangle$ ,  $|01\rangle$ ,  $|10\rangle$  and  $|11\rangle$ .

### 3.2.3 Measurement in bases other than the computational basis

If we try to represent state of a system  $|\psi\rangle = \alpha|0\rangle + \beta|1\rangle$  in a different bases,  $|a\rangle$  and  $|b\rangle$  which are *orthonormal* as  $\alpha'|a\rangle + \beta'|b\rangle$ . By postulates of quantum mechanics, if we try to measure w.r.t basis  $|a\rangle$ ,  $|b\rangle$  we get  $a$  with probability  $|\alpha'|^2$  and  $b$  with probability  $|\beta'|^2$ . We'll learn more about how efficient is it to measure w.r.t a basis in further study.

### 3.2.4 Quantum circuits

Similar to classical circuits, quantum circuits are made of wires, which don't necessarily represent physical wires, rather the passage of time or a physical photon moving in space. Let's first look at an example.

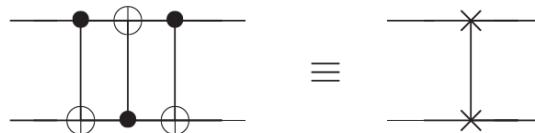


Figure 3.5: Circuit which swaps two qubits and an equivalent circuit since this is more useful and commonly used.

The above circuit swaps two qubits, which can be seen this way considering two qubits with states  $|a\rangle$  and  $|b\rangle$  as input,

$$|a, b\rangle \rightarrow |a, b \oplus a\rangle \quad (3.14)$$

$$\rightarrow |a \oplus (b \oplus a), b \oplus a\rangle = |b, b \oplus a\rangle \quad (3.15)$$

$$\rightarrow |b, (b \oplus a) \oplus b\rangle = |b, a\rangle \quad (3.16)$$

There are few features in classical circuits which aren't seen in quantum circuits. First one is *loops* are not allowed, i.e the circuit is *acyclic*. Second, classical circuit allows joining of wires,

an operation known as **FANIN**. Here, the bitwise OR operator of inputs is taken. However, information is lost and thus this is irreversible, thus not being allowed in quantum circuits. Thirdly, the inverse of this, where multiple copies of a single bit is made in a classical circuits, known as **FANOUT** is not allowed in quantum circuits. We can show that it's impossible to copy a qubit.

There's another important gate, which is in some sense generalisation of CNOT gate, *controlled-U* gate. which is represented by this

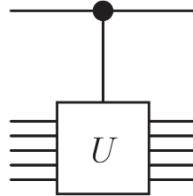


Figure 3.6: Controlled-U gate

This has similar logic, it's associated with a *unitary matrix*  $U$  and takes 1 qubit as the *control qubit*, and  $n$  qubits as *target qubits* as inputs. If the control qubit is set to 0, then nothing happens to target qubits. If control qubit is set to 1, then the gate  $U$  is applied to all the target qubits. In this sense, CNOT gate is the special case when  $n = 1$  and  $U = X$ , as shown



Figure 3.7: Two different representations of CNOT gate.

There's another important operation, known as measurement shown by the 'meter' symbol.

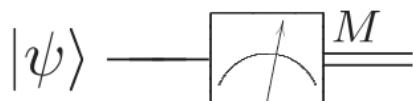


Figure 3.8: Quantum circuit symbol for measurement.

As expected, this outputs a probabilistic classical bit  $M$  (which is 0 with probability  $|\alpha|^2$  or 1 with probability  $|\beta|^2$ ) if a qubit with state  $|\psi\rangle = \alpha|0\rangle + \beta|1\rangle$  is given as input. Thus, the output is distinguished from a qubit by drawing a double-line wire.

### 3.2.5 Qubit copying circuit?

First we'll look at a classical example, suppose we have a bit  $x$  which we want to copy, we'll use a CNOT gate with  $x$  and a 'scratchpad bit' 0 as first and second inputs respectively, it'll copy  $x$  successfully as shown

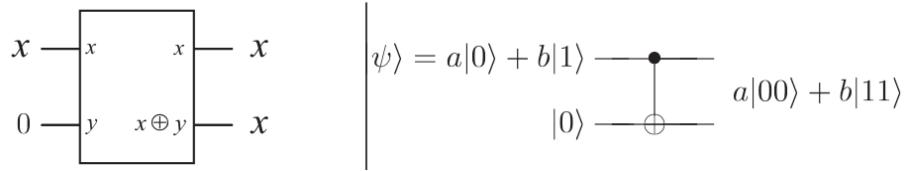


Figure 3.9: Trying to copy a classical bit (left) and a qubit (right)

But if we try copying a qubit with state  $|\psi\rangle = a|0\rangle + b|1\rangle$  in a similar manner using a CNOT gate, we can give  $|0\rangle$  as an input along with it as

$$[a|0\rangle + b|1\rangle]|0\rangle = a|00\rangle + b|10\rangle \quad (3.17)$$

applying the CNOT gate gives  $a|00\rangle + b|11\rangle$  as the output. We'd expect the output  $|\psi\rangle|\psi\rangle$  if we copied  $|\psi\rangle$  successfully. But that would mean the output should be

$$|\psi\rangle|\psi\rangle = a^2|00\rangle + ab|01\rangle + ab|10\rangle + b^2|11\rangle \quad (3.18)$$

But as we've seen above, it's only possible when  $ab = 0$ , i.e  $a = 0$  or  $b = 0$  i.e we want to copy either  $|0\rangle$  or  $|1\rangle$ . Hence we can't copy  $|\psi\rangle$ . It can be shown in general that it is impossible to copy a general qubit with unknown quantum state, if we construct a cloning gate, it can copy different state qubits  $|a\rangle$  and  $|b\rangle$  only if they are orthogonal. This is known as the *no-cloning theorem*.

There's another way to look at it. An unknown qubit contains some 'hidden' information that isn't directly accessible with a measurement. What's happening above is we get  $a|00\rangle + b|11\rangle$ , which if we measure we obtain either 0 or 1 using which the other qubit is completely determined. Thus, no extra information is gained. But if we were able to copy  $|\psi\rangle$ , then the other qubit contains some more information, contradicting our knowledge, thus a copy can't be created.

### 3.2.6 Example: Bell states

Let's take a look at a circuit which has a Hadamard gate on first input followed by a CNOT gate, as shown

In	Out
$ 00\rangle$	$( 00\rangle +  11\rangle)/\sqrt{2} \equiv  \beta_{00}\rangle$
$ 01\rangle$	$( 01\rangle +  10\rangle)/\sqrt{2} \equiv  \beta_{01}\rangle$
$ 10\rangle$	$( 00\rangle -  11\rangle)/\sqrt{2} \equiv  \beta_{10}\rangle$
$ 11\rangle$	$( 01\rangle -  10\rangle)/\sqrt{2} \equiv  \beta_{11}\rangle$

The quantum circuit on the right has two wires. The top wire starts with an H gate (represented by a box labeled 'H'), followed by a CNOT gate (control at top, target at bottom), and ends with the state  $|\beta_{xy}\rangle$ . The bottom wire starts with a CNOT gate (control at bottom, target at top), and ends with the state  $|\beta_{xy}\rangle$ .

Figure 3.10: Quantum circuit which creates Bell states, along with its quantum 'truth table'

As shown in the truth table, the Hadamard gate first transforms the first qubit into a superposition after which the CNOT gate inverts the target only when control is 1. The output

states are

$$|\beta_{00}\rangle = \frac{|00\rangle + |11\rangle}{\sqrt{2}} \quad (3.19)$$

$$|\beta_{01}\rangle = \frac{|01\rangle + |10\rangle}{\sqrt{2}} \quad (3.20)$$

$$|\beta_{10}\rangle = \frac{|00\rangle - |11\rangle}{\sqrt{2}} \quad (3.21)$$

$$|\beta_{11}\rangle = \frac{|01\rangle - |10\rangle}{\sqrt{2}} \quad (3.22)$$

These are known as the *Bell states* or *EPR states* or *EPR pairs* as their strange properties were found by Einstein, Podolsky and Rosen. Generalizing it,

$$|\beta_{xy}\rangle = \frac{|0y\rangle + (-1)^x |1\bar{y}\rangle}{\sqrt{2}} \quad (3.23)$$

where  $\bar{y}$  is the negation of  $y$ ,  $|\beta_{xy}\rangle$  denotes the output when the input given is  $|xy\rangle$ .

### 3.2.7 Example: quantum teleportation

It's a non-trivial, surprising and a fun thing! Quantum teleportation is a technique using which we can move quantum states around even in the absence of a quantum communication channel. Consider Alice and Bob, who've met long ago and have shared an EPR pair. Now Alice has a new qubit in state  $|\psi\rangle$  which she doesn't know and has to deliver to Bob. She can only send *classical* information.

Intuitively, it's pretty bad situation. It's not possible for her to get to know the state  $|\psi\rangle$  completely as she only has one copy of the qubit. Also if she knew the state, it's impossible to transfer this information to Bob, as  $|\psi\rangle$  takes values in *continuous* space, hence requiring infinite amount of time. Quantum teleportation provides a way for it.

In short, what happens is she measures this state  $|\psi\rangle$  along with the half of her EPR pair. She gets a result in 00, 01, 10 or 11 which she then sends to Bob. Using this information Bob can perform some operations and completely retrieve the state  $|\psi\rangle$  without even getting to know it. Shown using a quantum circuit

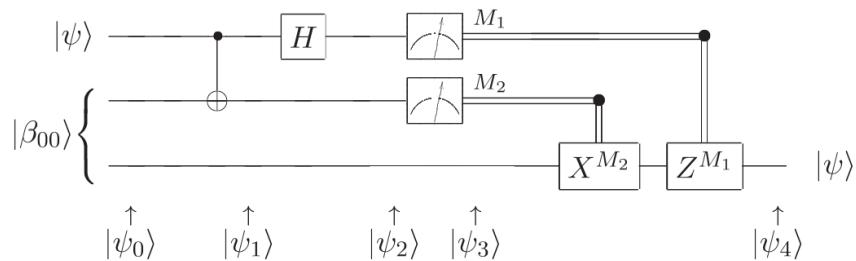


Figure 3.11: Quantum circuit showing the teleportation process

Suppose Alice and Bob are sharing  $|\beta_{00}\rangle$  and the  $|\psi\rangle$ , assuming second input is from

alice, overall input to the circuit is

$$|\psi_0\rangle = |\psi\rangle |\beta_{00}\rangle \quad (3.24)$$

$$= (\alpha|0\rangle + \beta|1\rangle) \left( \frac{|00\rangle + |11\rangle}{\sqrt{2}} \right) \quad (3.25)$$

$$= \frac{1}{\sqrt{2}} [\alpha|0\rangle|00\rangle + \alpha|0\rangle|11\rangle + \beta|1\rangle|00\rangle + \beta|1\rangle|11\rangle] \quad (3.26)$$

which when sent through CNOT gate on first two qubits gives  $|\psi_1\rangle$

$$|\psi_1\rangle = \frac{1}{\sqrt{2}} [\alpha|0\rangle(|00\rangle + |11\rangle) + \beta|1\rangle(|10\rangle + |01\rangle)] \quad (3.27)$$

The first qubit is then sent through Hadamard gate, which upon simplifying gives

$$|\psi_2\rangle = \frac{1}{2} [\alpha(|0\rangle + |1\rangle)(|00\rangle + |11\rangle) + \beta(|0\rangle - |1\rangle)(|10\rangle + |01\rangle)] \quad (3.28)$$

Now Alice measures the first two qubits of this whole system, giving outputs as  $M_1$  and  $M_2$  respectively. Based on what is measured as  $M_1 M_2$  we can know what is the final state of the qubit Bob has, which is

$$00 \mapsto |\psi_3(00)\rangle \equiv [\alpha|0\rangle + \beta|1\rangle] \quad (3.29)$$

$$01 \mapsto |\psi_3(01)\rangle \equiv [\alpha|1\rangle + \beta|0\rangle] \quad (3.30)$$

$$10 \mapsto |\psi_3(10)\rangle \equiv [\alpha|0\rangle - \beta|1\rangle] \quad (3.31)$$

$$11 \mapsto |\psi_3(11)\rangle \equiv [\alpha|1\rangle - \beta|0\rangle] \quad (3.32)$$

Now Alice sends these two classical bits to Bob, according to which he knows what kind of transformation to apply on his qubit to transform it to the qubit Alice had. We can verify that he has to apply the matrix  $Z^{M_1} X^{M_2}$  to get  $|\psi_4\rangle$  i.e  $|\psi\rangle$ . Note that the order of matrix multiplication is opposite to the time flow as shown in the circuit diagram.

As we had to transfer classical bits, which can't be done faster than the speed of light, quantum teleportation also can't be done faster than light and is limited by the speed of transfer of the classical bits as faster than light travel could possibly result in sending information backwards in time.

Also note that we haven't created a copy of  $|\psi\rangle$  which would contradict our *no-cloning theorem*. It's because the origin qubit we had in state  $|\psi\rangle$  changes into the state  $|0\rangle$  or  $|1\rangle$  depending on the value of  $M_1$  at the end of the process. Quantum teleportation shows that two classical bits can become a resource at least equal to one qubit information.

### 3.3 Quantum Algorithms

We'll explore what kind of problems a quantum computer can solve better than a classical computer. How a quantum computer can solve classical problems, we'll also look through well known quantum algorithms.

#### 3.3.1 Classical algorithms on a quantum computer

It turns out that we can simulate any classical circuit using a quantum circuit. It should be possible as physicists believe that everything is explainable using quantum mechanics. We'll

kind of prove it in the further discussion. Also, the reason we don't directly use quantum gates to simulate classical gates is that quantum gates are inherently *reversible*, whereas many classical gates such as NAND gates are inherently irreversible.

Any classical circuit can be replaced by a reversible elements known as the *Toffoli gate*. Toffoli gate is a reversible gate which takes in three input bits  $a, b, c$ . If we both the first two input bits (*control* bits)  $a, b$  are 1, it flips the third bit (*target* bit) else it doesn't affect it. It's equivalent to  $(a, b, c) \rightarrow (a, b, c \oplus ab)$  and is represented by this

Inputs			Outputs		
$a$	$b$	$c$	$a'$	$b'$	$c'$
0	0	0	0	0	0
0	0	1	0	0	1
0	1	0	0	1	0
0	1	1	0	1	1
1	0	0	1	0	0
1	0	1	1	0	1
1	1	0	1	1	1
1	1	1	1	1	0

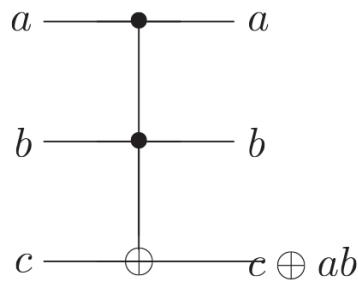


Figure 3.12: Left: Truth table for the Toffoli gate, Right: circuit representation of the Toffoli gate.

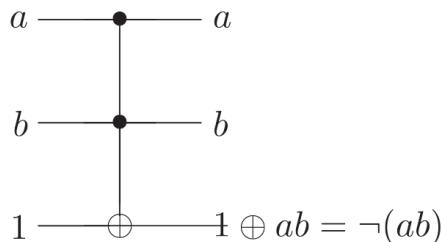


Figure 3.13: NAND gate simulated using toffoli gate.

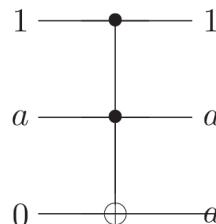


Figure 3.14: FANOUT gate simulated using Toffoli gate

Toffoli gate can also simulate NAND and FANOUT gates as shown above. Even though we've shown the classical implementation of toffoli gate above, it can be implemented as a quantum logic gate. It works in a similar way as the classical toffoli gate. For example flipping third qubit of  $|110\rangle$  to give  $|111\rangle$ . It can also be proved that this is a valid unitary quantum gate by constructing the corresponding 8 by 8 matrix,  $U$ . Thus, the quantum toffoli gate can simulate irreversible classical gates, ensuring that quantum computers are capable

of what classical computers can do. Not only a classical computer, a quantum computer can simulate a non-deterministic computer, i.e which can generate random bits. A simple way it can do it is preparing a state  $|0\rangle$  then applying the Hadamard gate to it, giving  $\frac{|0\rangle+|1\rangle}{\sqrt{2}}$  which upon measuring gives  $|0\rangle$  or  $|1\rangle$  with 50% probability, thus generating a random number and efficiently simulating a non-deterministic classical computer. A quantum computer is way more advantageous than this and can perform much more powerful functions as we'll discuss further.

### 3.3.2 Quantum parallelism

*Quantum parallelism* is a fundamental feature of quantum computers. They can evaluate a function  $f(x)$  simultaneously at many different values.

Suppose  $f : \{0, 1\} \rightarrow \{0, 1\}$  be a function with a one-bit domain and range. A way of computing this function using a quantum computer is to generate a transform using quantum logic gates which transforms a given state  $|x, y\rangle \rightarrow |x, y \oplus f(x)\rangle$ . When  $y = 0$  the state of the second qubit is just  $f(x)$ . We name this map as  $U_f$  as shown

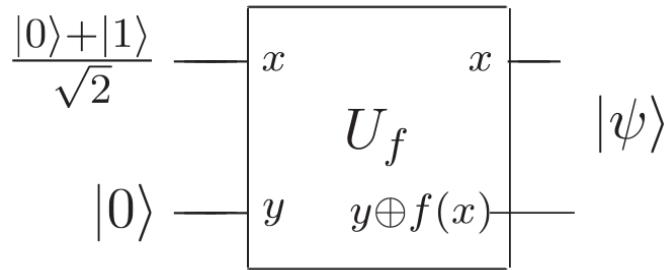


Figure 3.15: Quantum circuit for calculating  $f(0)$  and  $f(1)$  simultaneously.  $U_f$  is a quantum circuit which takes  $|x, y\rangle$  to  $|x, y \oplus f(x)\rangle$

Now if we want to compute both  $f(0)$  and  $f(1)$  simultaneously, we'll use the superposition state  $\frac{|0\rangle+|1\rangle}{\sqrt{2}}$ , which can be generated from  $|0\rangle$  using a Hadmard gate, and apply  $U_f$  on it. We then get

$$\frac{|0, f(0)\rangle + |1, f(1)\rangle}{\sqrt{2}} \quad (3.33)$$

which contains information about both  $f(1)$  and  $f(0)$ . Thus we've in some sense evaluated both  $f(0)$  and  $f(1)$  parallelly using a single circuit. In classical parallelism, we would have had to construct multiple circuits to evaluate the function at multiple points parallelly. We can generalize it to a function which takes *multiple* bits as inputs. First we'll apply *Hadamard transform* sometimes known as *Welsh-Hadamard transform* parallelly on  $n$  qubits. Circuit for  $n = 2$  is as shown

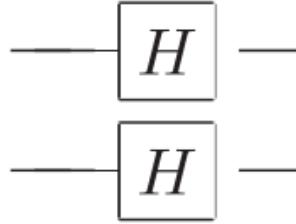


Figure 3.16: Two Hadmard gates ( $H^{\otimes 2}$ ) acting on two qubits parallelly.

For  $n = 2$  we get

$$\left( \frac{|0\rangle + |1\rangle}{\sqrt{2}} \right) \left( \frac{|0\rangle + |1\rangle}{\sqrt{2}} \right) = \frac{|00\rangle + |01\rangle + |10\rangle + |11\rangle}{2} \quad (3.34)$$

This operation can be represented by  $H^{\otimes 2}$ . For  $n$  qubits, as  $H^{\otimes n}$ . Performing Hadmard transform on  $n$  qubits all initialised to  $|0\rangle$  gives

$$\frac{1}{\sqrt{2^n}} \sum_x |x\rangle \quad (3.35)$$

Where  $x$  is all the possible combinations of  $n$  0's and 1's.

Now if we want to use quantum parallelism to compute  $f(x)$  which takes multi-bit values at multiple values, we first prepare a quantum system of  $n+1$  qubits with state  $|0\rangle^{\otimes n} |0\rangle$  which upon applying Hadmard transform on first  $n$  qubits gives 3.35. This followed by  $U_f$  produces

$$\frac{1}{\sqrt{2^n}} \sum_x |x\rangle |f(x)\rangle \quad (3.36)$$

Do note that we can still *measure* the value of  $f(x)$  only at one input. For quantum parallelism to be useful, we need some method to *extract* more information from superposition states like  $\sum_x |x\rangle |f(x)\rangle$ .

### 3.3.3 Deutsch's algorithm

Using this, we can get 'global information' about  $f(x)$ . It combines quantum parallelism with another quantum property known as *interference*. To see how this works, let's create our first qubit as superposition  $\frac{|0\rangle + |1\rangle}{\sqrt{2}}$  using Hadmard gate on  $|0\rangle$  and our second qubit as superposition  $\frac{|0\rangle - |1\rangle}{\sqrt{2}}$  using Hadmard gate on  $|1\rangle$ . Now let's use this circuit

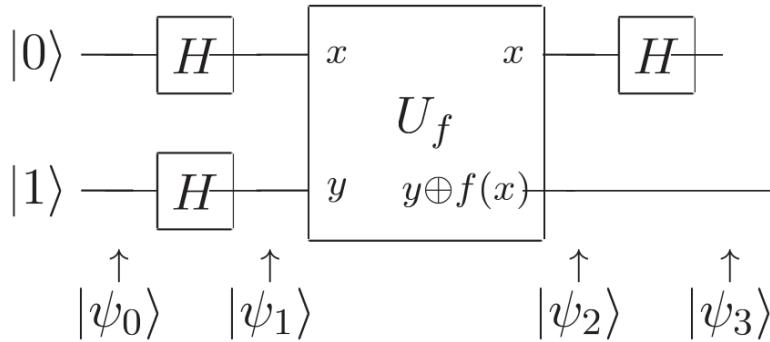


Figure 3.17: Quantum circuit implementing Deutsch's algorithm

Our input to this will be

$$|\psi_0\rangle = |0\rangle |1\rangle \quad (3.37)$$

When Hadmard gates are acted upon them, they become

$$|\psi_1\rangle = \left[ \frac{|0\rangle + |1\rangle}{\sqrt{2}} \right] \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] \quad (3.38)$$

With a little thought, it can be seen that applying  $U_f$  on  $|x\rangle (|0\rangle - |1\rangle)/\sqrt{2}$  gives  $(-1)^{f(x)} |x\rangle (|0\rangle - |1\rangle)/\sqrt{2}$ . Hence applying  $U_f$  to  $|\psi_1\rangle$  gives

$$|\psi_2\rangle = \begin{cases} \pm \left[ \frac{|0\rangle + |1\rangle}{\sqrt{2}} \right] \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] & \text{if } f(0) = f(1) \\ \pm \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] & \text{if } f(0) \neq f(1) \end{cases} \quad (3.39)$$

Applying the final Hadamard gate on this give  $|\psi_3\rangle$

$$|\psi_3\rangle = \begin{cases} \pm |0\rangle \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] & \text{if } f(0) = f(1) \\ \pm |1\rangle \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] & \text{if } f(0) \neq f(1) \end{cases} \quad (3.40)$$

This can be condensed more by observing that  $f(0) = f(1) \implies 0 = f(0) \oplus f(1)$  and  $f(0) \neq f(1) \implies 1 = f(0) \oplus f(1)$ . Thus

$$|\psi_3\rangle = \pm |f(0) \oplus f(1)\rangle \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] \quad (3.41)$$

Thus, we can measure  $f(0) \oplus f(1)$  by measuring the first qubit. This is really interesting since we found out  $f(0) \oplus f(1)$  which is a *global* property of  $f(x)$  by just evaluating  $f(x)$  once! Thus distinguishes a quantum parallelism from even a classical computer. Even a probabilistic classical computer could determine  $f(0)$  or  $f(1)$  with probability  $\frac{1}{2}$  each, a quantum computer evaluates both at once. We can get global information about function by cleverly choosing operation and gates so that the two evaluations *interfere* with each other to give us information like this, which is not possible in classical computers.

### 3.3.4 The Deutsch-Jozsa algorithm

This is a generalisation to the Deutsch's algorithm. Its application, known as *Deutsch's problem* is described as following. Suppose Alice and Bob are far apart, Alice is in Amsterdam and Bob is in Boston. Alice sends a number  $x$  out of  $\{0, \dots, 2^n - 1\}$  through a letter to Bob. Bob computes a function  $f(x)$  which yields the result of either 0 or 1 and mails it back to Alice. Alice needs to find out whether the function  $f(x)$  is *constant* or *balanced* i.e 0 for half the values and 1 for the other half by sending as many mails as she can. Using classical means we can figure out that she would need  $2^{n-1} + 1$  mails to be sent in the worst case.

We can offer a much better solution using qubits if we're sure that Bob is willing to calculate  $f(x)$  using a unitary transform  $U_f$  in a single interaction. To see how this works, Alice has a set of  $n$  qubits each prepared as  $|0\rangle$  in which she stores the query ( $\sim x$ ) and another qubit prepared as  $|1\rangle$  where she stores the answer ( $\sim f(x)$ ). She then applies Hadmard transform on all  $n + 1$  qubits after which she sends these to Bob. Bob first applies  $U_f$  on them and then Alice applies Hadmard transform on the first  $n$  qubits such that she gets

information about  $f(x)$  and thus the solution. We'll look through how this happens. Take a look at the circuit

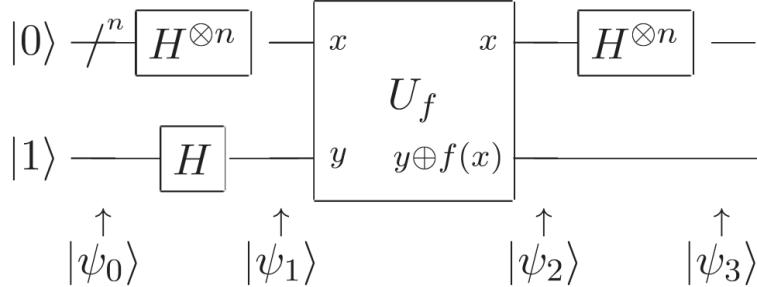


Figure 3.18: Implementation of Deutsch-Jozsa algorithm, the cross represents  $n$  qubits

So the input state is

$$|\psi_0\rangle = |0\rangle^{\otimes n} |1\rangle \quad (3.42)$$

When the Hadmard transform is applied, the query qubits turn into a superposition of  $|x\rangle$  where  $x$  is binary representation of all numbers in  $\{0, \dots, 2^n - 1\}$

$$|\psi_1\rangle = \sum_x \frac{|x\rangle}{\sqrt{2^n}} \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] \quad (3.43)$$

This is sent to Bob, he applies  $U_f$  to  $|\psi_1\rangle$  which results in  $|\psi_2\rangle$  using similar result as in 3.38

$$|\psi_2\rangle = \sum_x \frac{|x\rangle}{\sqrt{2^n}} (-1)^{f(x)} \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] \quad (3.44)$$

Now comes the tricky part when Alice applies Hadmard gate on first  $n$  qubits. Let's just see what happens when Hadamard transform is applied to  $|x\rangle$  which is just a binary string (representation of some number  $\in \{0, \dots, 2^n - 1\}$ ). We get another superposition of all possible basis states  $|z\rangle$  whose quotients are  $(-1)^{xz}$  which is bitwise multiplication/AND of  $x$  and  $z$  along with the normalization constant.

$$H|x\rangle = \sum_z \frac{(-1)^{xz}|z\rangle}{\sqrt{2^n}} \quad (3.45)$$

Thus the final output Alice gets is

$$|\psi_3\rangle = \sum_x \sum_z \frac{(-1)^{xz+f(x)}|z\rangle}{2^n} \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right] \quad (3.46)$$

Looking at first  $n$  qubits, coefficient of  $|0\rangle^{\otimes n}$  turns out to be just  $\sum_x \frac{(-1)^{f(x)}}{2^n}$ . Now note that if  $f(x)$  was constant  $\sum_x (-1)^{f(x)}$  would be  $\pm 2^n$ , which would mean coefficient of  $|0\rangle^{\otimes n}$  is  $\pm 1$ . Since norm of state of a system should be 1 coefficients of all other  $|z\rangle$  would be 0. Thus, Alice would only measure  $|0\rangle^{\otimes n}$ . Now in the case when  $f(x)$  is balanced, coefficient of  $|0\rangle^{\otimes n}$  would be 0. Hence, Alice would never get the measurement result as 0.

*Alice gets  $|0\rangle^{\otimes n}$  upon measuring  $|\psi_3\rangle \iff f(x) \text{ is constant}$*

Thus by measuring the query register Alice can get to know the answer in just one correspondence. Let's summarise this algorithm

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**Algorithm 1:** Deutsch-Jozsa
 

---

**Inputs :** A black box  $U_f$  which performs the transformation  
 $|x\rangle|y\rangle \rightarrow |x\rangle|y \oplus f(x)\rangle$ ,  $x \in \{0, \dots, 2^n - 1\}$  and  $f(x) \in \{0, 1\}$ . It's guaranteed that  $f(X)$  is either *balanced* or *constant*.

**Outputs :** 0 if and only if  $f$  is constant.

**Runtime:** One evaluation of  $U_f$ , always succeeds.

**Procedure**

```

 $|0\rangle^{\otimes n} |1\rangle$  // initialize state

 $\rightarrow \frac{1}{\sqrt{2^n}} \sum_{x=0}^{2^n-1} |x\rangle \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right]$  // create superposition using Hadmard gates

 $\rightarrow \sum_x (-1)^{f(x)} |x\rangle \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right]$  // calculate  $f$  using  $U_f$ 

 $\rightarrow \sum_z \sum_x \frac{(-1)^{xz+f(x)}}{2^n} |z\rangle \left[ \frac{|0\rangle - |1\rangle}{\sqrt{2}} \right]$  // perform Hadmard transform

 $\rightarrow z$  // measure to obtain output
  
```

---

This is a seed for quantum algorithms. But it has its own problems, such it has no real world application. It can be solved in a more realistic manner using a probabilistic computer.

### 3.3.5 Quantum algorithms summarized

Broadly speaking, quantum computers are good at solving three classes of problems. First one based upon quantum versions of Fourier transform, Deutsch-Jozsa algorithm, Shor's algorithm and discrete logarithm come under this class. Second class is quantum search algorithms. Third class of algorithms is quantum simulation, using which a quantum computer can simulate a quantum system. Let's look at each of these

#### Quantum algorithms based on Fourier transform

The discrete Fourier transform is described as transforming the set  $x_0, \dots, x_{N-1}$  of  $N$  complex numbers into the set of complex numbers  $y_0, \dots, y_{N-1}$  according to this

$$y_k \equiv \frac{1}{\sqrt{N}} \sum_{j=0}^{N-1} e^{\frac{2\pi i j k}{N}} x_j \quad (3.47)$$

It's widely known in science because turns a much complex problem into an easier one. There's a more generalized theory of fourier transforms which we won't be discussing. The Hadamard gates we've used in Deutsch-Jozsa algorithm and also the famous Shor's algorithm and discrete logarithm involve some kind of fourier transform. To go into the quantum realm, we can define a linear transform  $U$  on  $n$  qubits whose action on the basis states  $|k\rangle$ ,  $0 \leq k \leq 2^n - 1$  is defined as

$$|j\rangle \rightarrow \frac{1}{\sqrt{2^n}} \sum_{k=0}^{2^n-1} e^{\frac{2\pi i j k}{2^n}} |k\rangle \quad (3.48)$$

It can be checked that this transform is unitary, and can be realised as a quantum circuit. On a general qubit, this is

$$|\psi\rangle = \sum_{j=0}^{2^n-1} x_j |j\rangle \longrightarrow \frac{1}{\sqrt{2^n}} \sum_{k=0}^{2^n-1} \left[ \sum_{j=0}^{2^n-1} e^{\frac{2\pi i j k}{2^n}} x_j \right] |k\rangle = \frac{1}{\sqrt{2^n}} \sum_{k=0}^{2^n-1} y_k |k\rangle \quad (3.49)$$

Which can be seen as the vector form of fourier transform on the set  $x_j$  and  $N = 2^n$ .

A classical computer would perform the fourier transform in  $N \log N = n2^n$  steps, whereas a quantum computer would just take  $(\log N)^2 = n^2$  steps. Again the same problem arises that a lot of information is “hidden” and not accessible. When we try to measure it the state just collapses to either  $|0\rangle$  or  $|1\rangle$  making us loose our information. Hence, more cleverness is needed to harness the proper power of quantum computation. We’ll come back to more applications of Fourier Transform.

## Quantum search algorithms

A quantum computer can perform search algorithms, where we have to find an element from a collection of  $n$  random elements with a specific property in  $\sqrt{n}$  steps. Whereas a classical computer takes  $n$  steps. Even though this advantage is less than what we had in fourier transform’s case, this field has a wide variety of applications compared to the latter one.

## Quantum simulation

Simulating quantum system, as expected, is better done by quantum computers than classical ones. For a classical computers the number of complex numbers needed to describe a system grows *exponentially* whereas it grows *linearly* for a quantum computer i.e a quantum system having  $n$  components would require  $c^n$  bits of memory for a classical computer where  $c$  depends on the system and accuracy we need. Whereas a quantum computer would require only  $kn$  bits of memory. Again, the  $c^n$  bits of information isn’t accessible using the  $kn$  bit quantum computer. We’d still just be able to access  $kn$  bits of memory. How to extract more information is still partially understood. Still quantum simulation would be an important applications of QIC as we could understand about complex chemical molecules which generally occur in biology.

## The power of quantum computation

We’ll try to get some idea of how powerful quantum computers can be compared to classical computers, yet there is still a possibility they are not much powerful than a classical computer. *Computational complexity theory* deals with how difficult can it be to solve a problem. Most basic idea of it is of *complexity class* which is a set of problems all having some kind of similarity in how resource intensive they are.

Two most important complexity classes are **P** and **NP**. **P** is the set of problems which can be quickly solved using a classical computer. **NP** is the set of problems whose solutions can be quickly checked. For example, there’s no algorithm for factoring two integers that’s in **P** on a classical computer. But we can check if  $p$  is a factor of  $n$ , hence it’s an **NP** problem.

It’s clear that **P** is a subset of **NP** since the ability to solve a problem implies the ability to check potential solutions. But there are no known problems that are “surely” not in **P** but

are in **NP**. This is a famous open problem in computer science

$$\mathbf{P} \stackrel{?}{\neq} \mathbf{NP} \quad (3.50)$$

There's an important subclass of **NP** problems, known as **NP**-complete problems. Many hard and highly important problems are known to be **NP**-complete. In some sense, **NP**-complete problems are “atleast” as hard as **NP** problems. Any algorithm solving **NP**-complete problems can be adapted to solve **NP** hard problems, with a little overhead. If  $\mathbf{P} \neq \mathbf{NP}$ , then there would be no **NP**-complete problem efficiently solvable by a classical computer.

It is not known whether quantum computers can solve all **NP** problems. They can solve few problems like factoring which aren't surely known to be in **NP** but are believed to be in. There's an interesting result that a simple variant of quantum parallelism can't solve all **NP** problems. But still, this does not rule out that there might be some deeper structure which would let us solve all **NP** problems on a quantum computer efficiently.

Another important complexity class **PSPACE** is the set of problems which are efficient on space but not necessarily on time. Even though *not proved*, **PSPACE** is believed to be bigger than both **P** and **NP**. Finally, **BPP** is the set of problems that can be solved using randomized algorithms with some bound probability error (say 1/4). They're more believed to be the class that can be solved by a classical computer. Even though **P** is more studied.

**BQP** is the class of all computational problems that can be solved efficiently on a quantum computer with a bound probability error. **BQP** is knownn surely to be bigger than **P** but not bigger than **PSPACE**. It's not exactly known where it fits between **P**, **NP**, **PSPACE**. This gives an important fact that if *quantum computers are strictly powerful than classical computers* then **PSPACE** is strictly bigger than **P**. The latter part has been really hard to prove by computer scientists, lowering the hope that quantum computers are more powerful than classical computers.

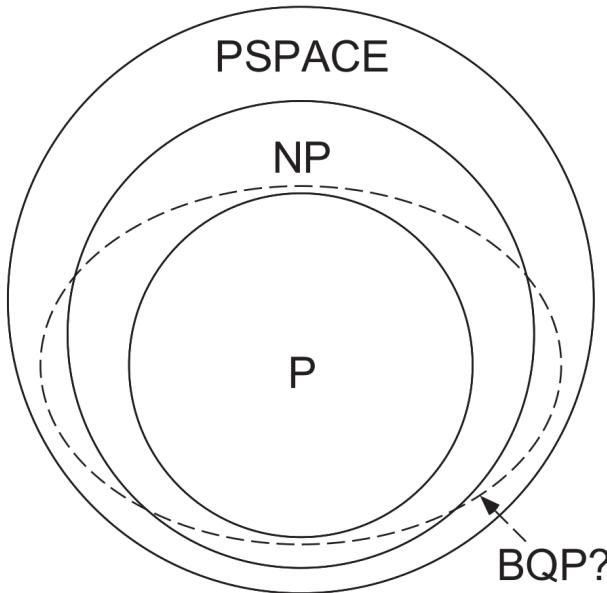


Figure 3.19: We surely know that quantum computers can solve all **P** problems quickly and also that they cannot solve any problem outside **PSPACE** but we don't know where quantum computers fit between **P** and **PSPACE**. We don't even know whether **P** is equal to **PSPACE**.

We can already see that the *theory* and notions of quantum computing pose a challenge to traditional computational methods. Another important challenge is that quantum com-

puting is believed to be *experimentally realizable*, primarily because nature works according to quantum mechanics.

## 3.4 Experimental quantum information processing

We'll look into experimentally proving that quantum computing can be done. Let's start with the "Stern-Gerlach" experiment which provides evidence to the existence of qubits in nature.

### 3.4.1 The Stern-Gerlach experiment

In the original experiment performed in 1921-22, hot 'silver atoms' were beamed from an oven through a magnetic field which deflected it. We'll consider the hydrogen atom version performed in 1927 as it's much simpler. A hydrogen atom consists of a proton and an electron spinning around which makes up the atomic dipole moment. By constructing the apparatus properly, it's possible to make the atom deflect by an amount dependent on the  $\hat{z}$  component of atomic dipole moment. It's naturally expected that the atomic dipole moments would be randomly oriented and we'd get a continuous distribution. But the atoms were deflecting at specific angles. This was then explained by the *quantization* principle, which was trending at that time. But quantum mechanics had predicted that the net atomic dipole of an atom must be zero, which is classically strange though, but there were two beams seen, one deflected up and the other deflected down.

This puzzle was explained by a new quantity called *spin* which every hydrogen atom is supposed to be associated with. This spin is posited to make *extra contribution* to the magnetic dipole moment.

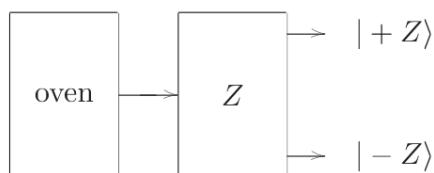


Figure 3.20: Original Stern-Gerlach setup.

We can assume that the spins  $|+Z\rangle$ ,  $| - Z \rangle$  leads to upper and lower deflection respectively. Now if the apparatus modified that the upper beam is then passed through another apparatus *tipped sideways* so that it deflects according to  $\hat{x}$  axis.

## **Part III**

### **Week 3**

# Chapter 4

## Quantum noise and quantum operations

We'll be discussing about *open systems*. These are unlike *closed systems*, which don't interact with outside environment. These can suffer with unwanted interactions with the outside environment. These unwanted interactions show up as *noise* in quantum information systems. We'll describe *quantum operations formalism*, a powerful set of tools enabling us to describe *open* quantum systems and their behaviour. These can describe not only systems weakly coupled with the environment, but also the system which are strongly coupled with the environment. They can also describe closed systems which are suddenly opened and are subject to measurement. They are good at describing *discrete* state changes, i.e changes of state from  $\rho$  to  $\rho'$  without explicit mention of passage of time.

### 4.1 Classical noise and Markov processes

To understand about noise, let's consider a simple situation where a hard drive of a computer stores a single bit, which can be 0 or 1. Due to external stray magnetic fields, this bit is subject to change after a long time. To quantify it, let  $p$  be the probability that the bit will flip, then  $1 - p$  will be the probability that the bit remains the same.

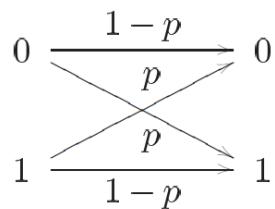


Figure 4.1: A bit flip might occur with probability  $p$

To figure out this probability  $p$ , we need to understand two things (a) How the magnetic fields are distributed in the environment. Assuming the owner isn't crazy and placing a magnet near the disk, we can sample the magnetic field near the disk similar to the one the disk is in. (b) How this environment's magnetic field interacts with the disk. This can be done by the well established *Maxwell's equations*.

Suppose  $p_0, p_1$  be the probabilities that the bit is in state 0 and 1 initially. Let  $q_0, q_1$  be the corresponding probabilities after the noise has occurred and  $X, Y$  be the initial and final

states respectively then

$$P(Y = y) = \sum_x P(Y = y|X = x)P(X = x) \quad (4.1)$$

where the conditional probability  $P(Y = y|X = x)$  is known as the *transition probability*. Thus

$$\begin{bmatrix} q_0 \\ q_1 \end{bmatrix} = \begin{bmatrix} 1-p & p \\ p & 1-p \end{bmatrix} \begin{bmatrix} p_0 \\ p_1 \end{bmatrix} \quad (4.2)$$

Suppose we make a quantum circuit made of two faulty NOT gates. It takes an input state  $X$  converts it into an intermediate bit  $Y$  which is finally converted to  $Z$ . If we make a physically reasonable assumption that the noise produced at first NOT gate independent of the noise produced at the second NOT gate, it results in a stochastic process  $X \rightarrow Y \rightarrow Z$  of a special type known as *Markov process*. Often in multi-stage processes, it's a good assumption to use Markov processes.

For a single stage process, output probabilities  $\vec{q}$  are related to the input probabilities  $\vec{p}$  according to

$$\vec{q} = E\vec{p} \quad (4.3)$$

$E$  is a matrix of transition probabilities known as the *Evolution matrix*. Thus the final state of system is linearly related to its initial state. This linearity is echoed in quantum noise's description, where probability distributions are replaced by density operators. In addition to this,  $E$  follows two more rules (a) All of its entries must be positive, so that  $E\vec{p}$  is also positive and a valid probability distribution. This is known as *positivity*. (b) Sum of the entries in each column should be 1. This is known as *completeness*.

## 4.2 Quantum operations

### 4.2.1 Overview

Quantum operations formalism is a general tool for describing evolution of quantum systems in a wide variety of circumstances, include stochastic changes, similar to how Markov processes describing classical stochastic changes. Similar to how classical bits are represented using probability distributions, here we'll use *density operators*. The transformation looks like

$$\rho' = \mathcal{E}(\rho) \quad (4.4)$$

Here,  $\rho$  is the initial density operator describing the state of the system. Similarly,  $\rho'$  is the final state upto a normalization factor. A few known transforms  $\mathcal{E}$  are unitary transformations and measurements  $\mathcal{E}(\rho) = U\rho U^\dagger$  and  $\mathcal{E}_m(\rho) = M_m \rho M_m^\dagger$ .

In our further study, we're going to look at quantum operations *three* separate ways. The first way is based on the idea of studying dynamics as the result of interaction between a system and environment. This method is *concrete* and easy to relate to real world. But it is mathematically inconvenient. Second way is equivalent to the first one but is mathematically convenient using a powerful mathematical representation for quantum operations known as the *operator-sum representation*. Third way is equivalent to the first two, it's via a set of *physically motivated axioms* that we'd expect these maps to satisfy. Its advantage is that it shows that quantum dynamics will be described by quantum operations in a wide range of circumstances. But this is neither concrete nor mathematically convenient.

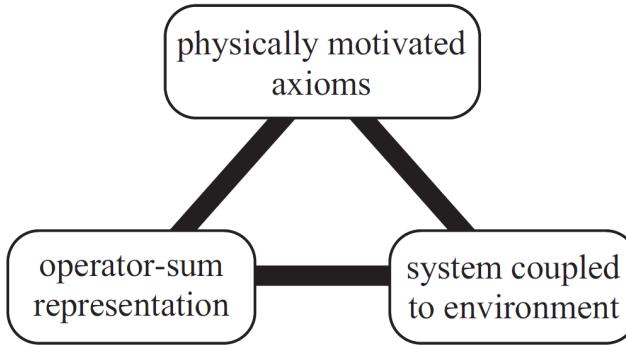


Figure 4.2: Three ways we're going to look at quantum operations, each provide their own advantages

### 4.2.2 Environment and quantum operations

By the postulate of quantum mechanics, dynamics of a *closed* system is described by just a unitary transform which can be thought of as a box. For dynamics of an *open* system, the system can be thought of as a composite closed system of the system we're considering, *principal system* and the *environment*.

Suppose the system is in state  $\rho$ , then the final state would be  $\mathcal{E}(\rho)$  which need not be a unitary transform. We *assume* that system and environment are in a product state  $\rho \otimes \rho_{\text{env}}$ . Evolution of this closed system is described by a unitary transform  $U$  and we use partial trace to find final state of the system  $\mathcal{E}(\rho)$

$$\mathcal{E}(\rho) = \text{tr}_{\text{env}} [U \rho \otimes \rho_{\text{env}} U^\dagger] \quad (4.5)$$

This is the first *definition* of quantum operations. If  $U$  doesn't depend on the environment, then final state is simply  $\tilde{U} \rho \tilde{U}^\dagger$  where  $\tilde{U}$  is just part of transform on our principal state.

We assumed that environment and principal system are in a product state, but this might not be the case always. In cases of practical interest, this assumption is reasonable. The experimentalist while preparing an open state clears all the *correlations* between the state and environment.

You might think the environment would have near-infinite dimensional Hilbert space, but it turns out that for principal system with  $d$ -dimensional Hilbert space, it's enough to model the environment with no more than  $d^2$  dimensions.

### 4.2.3 Operator sum representation

How an open system evolves can be understood with an operator acting on the Hilbert space of principal system. This representation of quantum operations is known as *Operator sum representation*.

Let this environment be in a pure state  $\rho_{\text{env}} = |e_0\rangle\langle e_0|$ , which even if it's not, can be made pure. Let  $|e_k\rangle$  be an orthonormal basis for environment. The operation can be represented as

$$\mathcal{E}(\rho) = \sum_k \langle e_k | U [\rho \otimes |e_0\rangle\langle e_0|] U^\dagger |e_k\rangle \quad (4.6)$$

$$= \sum_k E_k \rho E_k^\dagger \quad (4.7)$$

where  $E_k = \langle e_k | U | e_0 \rangle$  is an operator acting on principal state known as an *operation element*. These have a nice property known as *completeness relation* coming from the fact that  $\text{tr}(\mathcal{E}(\phi)) = 1$  since that's just another valid state

$$1 = \text{tr}(\mathcal{E}(\rho)) \quad (4.8)$$

$$= \text{tr} \left( \sum_k E_k \rho E_k^\dagger \right) \quad (4.9)$$

$$= \text{tr} \left( \sum_k E_k E_k^\dagger \rho \right) \quad (4.10)$$

As it's true for all  $\rho$ ,

$$\sum_k E_k E_k^\dagger = I \quad (4.11)$$

This is satisfied by operations which are *trace-preserving*. We'll see later non trace preserving operations are also equivalent to this.

This representation gives us *intrinsic* means to characterize dynamics of the principal system. We don't need to know how the environment works, just  $E_k$  are sufficient. Next up we'll see it's physical interpretation, making a physical model from this representation and also how to define this representation when our assumptions fail.

### Physical interpretation of operator-sum representation

Suppose a unitary transform  $U$  is applied followed by a measurement of the environment, in basis  $|e_k\rangle$ . The measurement only affects the environment, thus state of the system if the output is  $k$  is

$$\rho_k \propto \text{tr}_E (|e_k\rangle \langle e_k| U (\rho \otimes |e_0\rangle \langle e_0|) U^\dagger |e_k\rangle \langle e_k|) \quad (4.12)$$

$$= \langle e_k | U (\rho \otimes |e_0\rangle \langle e_0|) U^\dagger |e_k\rangle \quad (4.13)$$

$$= E_k \rho E_k^\dagger \quad (4.14)$$

which upon normalizing gives,

$$\rho_k = \frac{E_k \rho E_k^\dagger}{\text{tr}(E_k \rho E_k^\dagger)} \quad (4.15)$$

with probability

$$\text{tr}(E_k \rho E_k^\dagger) \quad (4.16)$$

thus the final output is

$$\mathcal{E}(\rho) = \sum_k E_k \rho E_k^\dagger \quad (4.17)$$

Thus we can see that a quantum operation changes the input randomly into a state  $\frac{E_k \rho E_k^\dagger}{\text{tr}(E_k \rho E_k^\dagger)}$  with probability  $\text{tr}(E_k \rho E_k^\dagger)$  which is very similar to a noisy channel.

### Measurement and operator-sum representation

We can have non trace preserving operations if we allow *measurement* of the principal state-environment system after the operator  $U$  is applied. Suppose the principal system is represented by  $Q$  and the environment by  $E$ , both are initially independent, environment is in the state  $\sigma$ . Then

$$\rho^{QE} = \rho \otimes \sigma \quad (4.18)$$

Now a unitary transform  $U$  is applied, then a projective measurement is made using  $P_m$ . We assume no measurement is made is same as the outcome is  $m = 0$  and  $P_0 \equiv I$ . The final state of  $QE$  is

$$\frac{P_m U(\rho \otimes \sigma) U^\dagger P_m}{\text{tr}(P_m U(\rho \otimes \sigma) U^\dagger P_m)} \quad (4.19)$$

State of only the principal system  $Q$  is

$$\frac{\text{tr}_E(P_m U(\rho \otimes \sigma) U^\dagger P_m)}{\text{tr}(P_m U(\rho \otimes \sigma) U^\dagger P_m)} \quad (4.20)$$

We can define a map  $\mathcal{E}_m(\rho)$  as

$$\mathcal{E}_m(\rho) = \text{tr}_E(P_m U(\rho \otimes \sigma) U^\dagger P_m) \quad (4.21)$$

therefore the final state of  $Q$  is  $\frac{\mathcal{E}_m(\rho)}{\text{tr}(\mathcal{E}_m(\rho))}$  and  $\text{tr}(\mathcal{E}_m(\rho))$  is the probability of measurement outcome  $m$ . Suppose  $|e_k\rangle$  is an orthonormal basis of  $E$  and  $\sigma = \sum_j q_j |j\rangle \langle j|$  is an ensemble of initial state of environment then we observe that

$$\mathcal{E}_m(\rho) = \sum_{jk} q_j \text{tr}_E(|e_k\rangle \langle e_k| P_m U(\rho \otimes \sigma) U^\dagger P_m |e_k\rangle \langle e_k|) \quad (4.22)$$

$$= \sum_{jk} E_{jk} \rho E_{jk}^\dagger \quad (4.23)$$

where

$$E_{jk} \equiv \sqrt{q_j} \langle e_j | U | j \rangle \quad (4.24)$$

using which we can calculate the operators for the operator sum representation if  $\sigma$  of  $E$  is known. Even though we don't make projective measurements, respective measurement probabilities are  $\text{tr}(\mathcal{E}_m(\rho))$

### System-environment models for any operator-sum representation

Now a natural question is for a given set of operators  $\{E_k\}$  is it possible to find a reasonable *model environment system and dynamics* which give rise to the quantum operation described by these operators. It turns out that for a set of trace preserving or non trace preserving quantum operation  $\mathcal{E}$ , there exists a model environment  $E$  starting with a pure state  $|e_0\rangle$ , and model dynamics specified by unitary operator  $U$  and projector  $P$  onto  $E$  such that

$$\mathcal{E}(\rho) = \text{tr}_E(P U(\rho \otimes |e_0\rangle \langle e_0|) U^\dagger P) \quad (4.25)$$

#### 4.2.4 Axiomatic approach to quantum operations

We'll try to write down physically motivated axioms which quantum operations obey. We'll start over from scratch using these axioms. We'll then prove that a map  $\mathcal{E}$  satisfies these axioms *if and only if* it has an operator sum representation.

We define a quantum operation  $\mathcal{E}$  as a map from the set of density operators of the input space  $Q_1$  to the set of density operators for the output space  $Q_2$ , which satisfy the following (for convenience we'd take  $Q_1 = Q_2 = Q$ ):

1.  $\text{tr}(\mathcal{E}(\rho))$  denotes the probability that the map described by  $\mathcal{E}$  occurs. Thus  $0 \leq \text{tr}(\mathcal{E}(\rho)) \leq 1$ .
2.  $\mathcal{E}$  is a *convex-linear map* on the set of density matrices i.e

$$\mathcal{E}\left(\sum_i p_i \rho_i\right) = \sum_i p_i \mathcal{E}(\rho_i). \quad (4.26)$$

3.  $\mathcal{E}$  is a *completely-positive map* i.e  $\mathcal{E}(A)$  is positive for any positive  $A$ . Also, if we introduce an extra system  $\mathcal{R}$ , then  $\mathcal{I} \otimes \mathcal{E}(A)$  is positive for any positive  $A$  in  $\mathcal{R}Q_1$ , where  $\mathcal{I}$  is the identity map on  $\mathcal{R}$ .

The following theorem binds the axiomatic approach and the operator-sum representation

**Theorem 8.** *The map  $\mathcal{E}$  satisfies the above axioms if and only if*

$$\mathcal{E}(\rho) = \sum_k E_k \rho E_k^\dagger \quad (4.27)$$

for some set of operators  $\{E_k\}$  mapping the input Hilbert space to the output Hilbert space and satisfying  $\sum_k E_k E_k^\dagger \leq I$

#### Freedom in the operator-sum representation

A quanum operation can be described by two different set of operators. For example take  $\mathcal{E}(\rho) = \sum_k E_k \rho E_k^\dagger$  and  $\mathcal{F}(\rho) = \sum_k F_k \rho F_k^\dagger$  and

$$E_1 = \frac{I}{\sqrt{2}} = \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \quad E_2 = \frac{Z}{\sqrt{2}} = \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix} \quad (4.28)$$

$$F_1 = |0\rangle \langle 0| = \begin{bmatrix} 1 & 0 \\ 0 & 0 \end{bmatrix} \quad F_2 = |1\rangle \langle 1| = \begin{bmatrix} 0 & 0 \\ 0 & 1 \end{bmatrix} \quad (4.29)$$

Understanding which sets of operations gives rise to same quantum operations is important for two reasons. (a) Understanding the freedom in representation gives us more insight into how different physical processes give rise to same dynamics (b) Understanding freedom in operator-sum representation is crucial to a good understanding of quantum error-correction. This is described usin the following theorem

**Theorem 9 (Unitary freedom in the operator-sum representation).** *Suppose  $\{E_1, \dots, E_m\}$  and  $\{F_1, \dots, F_n\}$  are operation elements giving rise to quantum operations  $\mathcal{E}$  and  $\mathcal{F}$  respectively. We may ensure  $m = n$  by appending 0 operators to the end of shorter operation elements' list. Then  $\mathcal{E} = \mathcal{F}$  if and only if there exist complex numbers  $u_{ij}$  such that  $E_i = \sum_j u_{ij} F_j$  and  $u_{ij}$  is a unitary matrix.*

This is similar to the result when two states are denoted by the same density operator. Using this theorem we can also answer the question, what is the maximum number of dimensions of an environment needed to mock up a quantum operation as described by this theorem.

**Theorem 10.** *All quantum operations  $\mathcal{E}$  on a principal system of Hilbert space dimension  $d$  can be generated by an operator sum representation containing atmost  $d^2$  elements, i.e*

$$\mathcal{E}(\rho) = \sum_{k=1}^M E_k \rho E_k^\dagger \quad (4.30)$$

where  $1 \leq M \leq d^2$ .

## 4.3 Examples of quantum noise and quantum operations

### 4.3.1 Trace and partial trace

Quantum operations formalism can be used to describe a measurement, it's outcome and the state of the system after the measurement. The simplest such operation is related to the map  $\rho \rightarrow \text{tr}(\rho)$ . Let  $H_Q$  be the input Hilbert space with orthogonal basis  $|1\rangle, \dots, |d\rangle$  and  $H'_Q$  is a one dimensional Hilbert space, spanned by  $|0\rangle$ , define  $\mathcal{E}(\rho)$  as

$$\mathcal{E}(\rho) \equiv \sum_{i=1}^d |0\rangle \langle i| \rho |i\rangle \langle 0| \quad (4.31)$$

Do note that  $\mathcal{E}(\rho) = \text{tr}(\rho) |0\rangle \langle 0|$ , thus this quantum operation is identical to the trace operation, ignoring the factor of  $|0\rangle \langle 0|$ .

Another interesting thing comes when we have a joint system  $QR$  and we want to trace out  $R$ . Let  $|j\rangle$  be the basis of  $R$ , let's define a linear operator  $E_i : H_{QR} \rightarrow H_Q$  as

$$E_i \left( \sum_j \lambda_j |q_j\rangle |j\rangle \right) = \lambda_i |q_i\rangle \quad (4.32)$$

when  $|q_j\rangle$  are arbitrary vectors in  $Q$  and  $\lambda_j$  are just constants. Then we define a quantum operation as

$$\mathcal{E}(\rho) \equiv \sum_k E_k \rho E_k^\dagger \quad (4.33)$$

We also notice that

$$\mathcal{E}(\rho \otimes |j\rangle \langle j'|) = \rho \delta_{j,j'} = \rho \text{tr}(|j\rangle \langle j'|) = \text{tr}_R(\rho \otimes |j\rangle \langle j'|). \quad (4.34)$$

where  $|j\rangle, |j'\rangle$  are members of orthonormal basis of  $R$ . Thus, by linearity of  $\mathcal{E}$  and  $\text{tr}_R$  we have  $\mathcal{E} = \text{tr}_R$

### 4.3.2 Geometric picture of single qubit quantum operations

We've seen before that a single qubit's density operator can be represented by

$$\rho = \frac{I + \vec{r} \cdot \vec{\sigma}}{2} \quad (4.35)$$

where  $\vec{r}$  is a three dimensional real unit vector and  $\vec{\sigma}$  is a vector of pauli matrices. Thus

$$\rho = \frac{1}{2} \begin{bmatrix} 1 + r_z & r_x - ir_y \\ r_x + ir_y & 1 - r_z \end{bmatrix} \quad (4.36)$$

Using this representation, it can be shown that an arbitrary trace-preserving quantum operation is equivalent to the map

$$\vec{r} \longrightarrow \vec{r}' = M\vec{r} + \vec{c} \quad (4.37)$$

where  $M$  is a  $3 \times 3$  real matrix, and  $\vec{c}$  is a constant vector. This is an *affine map*, mapping Bloch sphere into itself. To see this, suppose the operators  $E_i$  generating  $\mathcal{E}$  are written in the form

$$E_i = \alpha_i I + \sum_{k=1}^3 a_i k \sigma_k \quad (4.38)$$

Then it can be checked that

$$M_{jk} = \sum_l \left[ a_{lj} a_{lk}^* + a_{lj}^* a_{lk} + \left( |\alpha_l|^2 - \sum_p a_{lp} a_{lp}^* \right) \delta_{jk} + i \sum_p \epsilon_{jkl} (\alpha_l a_{lp}^* - \alpha_l^* a_{lp}) \right] \quad (4.39)$$

$$c_k = 2i \sum_l \sum_{jp} \epsilon_{jpk} a_{lj} a_{lp}^* \quad (4.40)$$

The affine map can be better understood by the polar decomposition of  $M$  into  $M = U|M|$ , where  $U$  is unitary. Since  $M$  is real,  $|M|$  is real and hermitian, i.e a symmetric matrix. Thus, as  $M$  is real, we may assume  $U$  is real, and an orthogonal matrix i.e  $U^T U = I$ . Thus, we can write

$$M = OS \quad (4.41)$$

where determinant of  $O$  is 1 and it's an orthogonal matrix describing rotation.  $S$  is a real symmetric matrix. Thus the above equation can be understood as the deformation of the bloch sphere about axis as specified by  $S$  followed by a rotation defined by  $O$ , then a displacement due to  $\vec{c}$ .

### 4.3.3 Bit flip and phase flip channels

Suppose we take a single qubit and pass it through a noisy *bit-flip* channel, which flips the state of a qubit from  $|0\rangle$  to  $|1\rangle$  with probability  $1 - p$ . It has operation elements

$$E_0 = \sqrt{p}I = \sqrt{p} \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \quad E_1 = \sqrt{1-p}X = \sqrt{1-p} \begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix} \quad (4.42)$$

This affects the bloch vector in the following way:

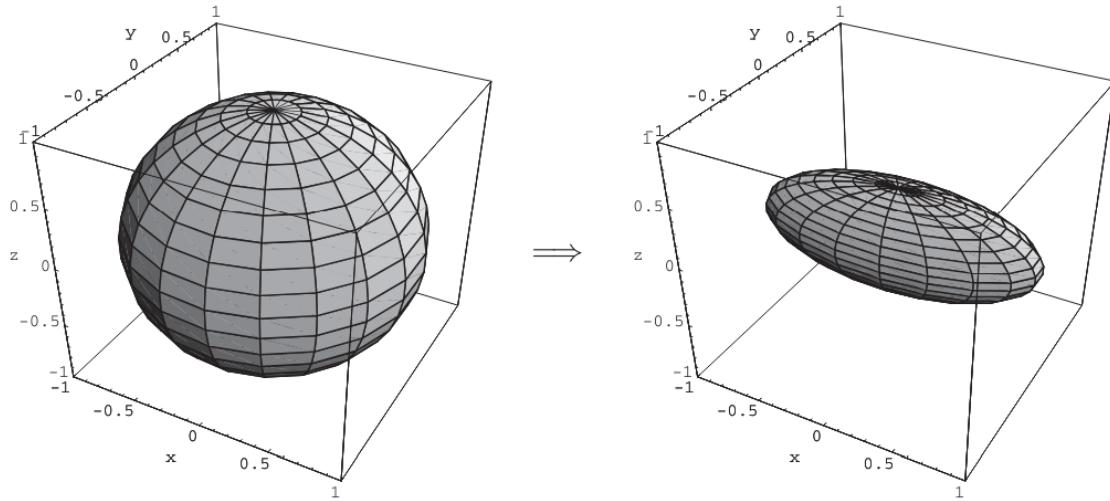


Figure 4.3: Bit flip channel. The operation doesn't affect the  $x$  component, but shrinks the  $y$  and  $z$  component by a factor of  $1 - 2p$ . Here  $p = 0.3$ .

It's easy to verify that  $\text{tr}(\rho^2) = \frac{1+|\vec{r}|^2}{2}$ . Thus, this operation can never increase the length of Bloch vector i.e can never increase  $\text{tr}(\rho^2)$ . Thus the purity of  $\rho$  is always decreased by this. Thus Bloch sphere provides an intuitive way of understanding quantum operations.

Similarly if we consider the *phase-flip* channel which has operation elements

$$E_0 = \sqrt{p}I = \sqrt{p} \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \quad E_1 = \sqrt{1-p}Z = \sqrt{1-p} \begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix} \quad (4.43)$$

we see that the following transformation is happening to the Bloch sphere.

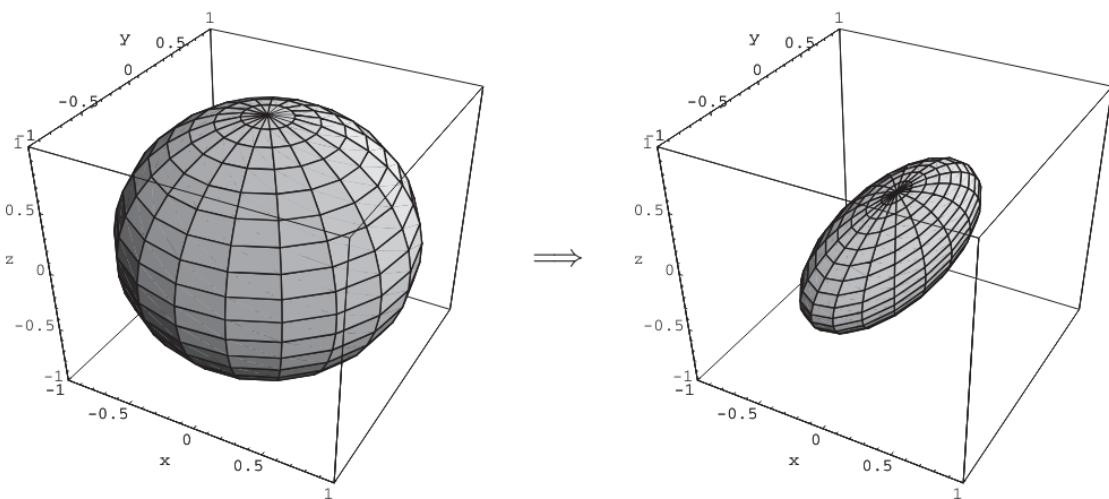


Figure 4.4: Phase flip channel. As we can see, it does nothing to the  $z$  component, but shrinks  $x$  and  $y$  component by a factor of  $1 - 2p$ . Here  $p=0.3$ .

As a special case when  $p = 0.5$ , using the freedom in the operator-sum representation, this operation can be written as

$$\rho \longrightarrow \mathcal{E}(\rho) = P_0\rho P_0 + P_1\rho P_1 \quad (4.44)$$

which is nothing but the measurement with respect to the basis  $|0\rangle, |1\rangle$ , the result unknown. The corresponding map on Bloch sphere here is

$$(r_x, r_y, r_z) \longrightarrow (0, 0, r_z) \quad (4.45)$$

i.e the Bloch vector is projected along  $z$  axis loosing  $x$  and  $y$  components.

The *bit-phase flip* channel has the operation elements

$$E_0 = \sqrt{p}I = \sqrt{p} \begin{bmatrix} 1 & 0 \\ 0 & 1 \end{bmatrix} \quad E_1 = \sqrt{1-p}Y = \sqrt{p} \begin{bmatrix} 0 & -i \\ i & 0 \end{bmatrix} \quad (4.46)$$

when this operation is done, the  $y$  component of the Bloch sphere remain intact. Whereas  $x, z$  components are shrunked by a factor of  $1 - 2p$  as shown

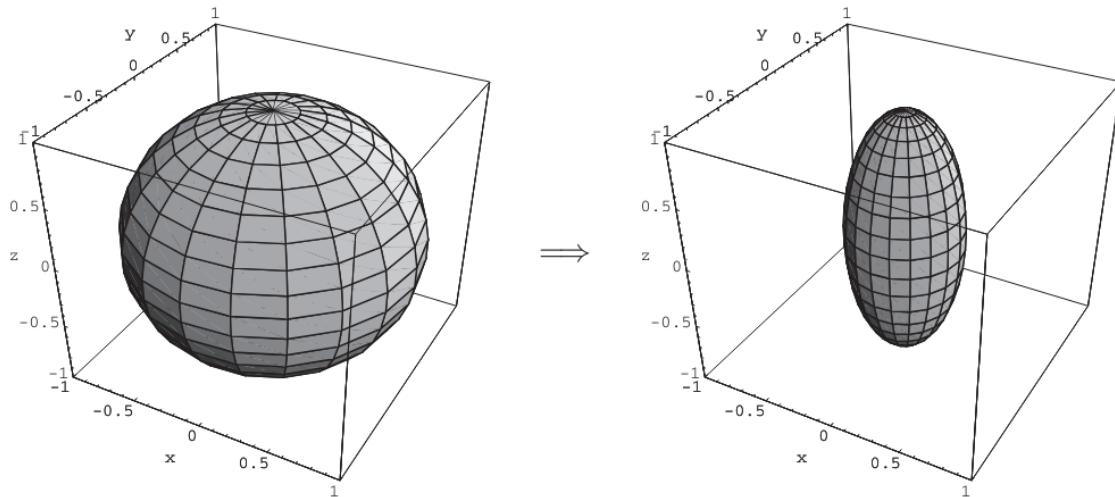


Figure 4.5: Bit-phase flip channel. Here  $p=0.3$ , it can be seen that  $y$  axis remains intact while  $x$  and  $z$  axes are shrunked by a factor of  $1 - 2p$ .

#### 4.3.4 Depolarizing channel

It's an important type of quantum noise channel. We take a single qubit, and with probability  $p$  the qubit is *depolarized* i.e replaced by a completely mixed state  $\frac{I}{2}$  or left untouched with probability  $1 - p$ . Thus the state of the qubit after noise is

$$\mathcal{E}(\rho) = \frac{pI}{2} + (1 - p)\rho \quad (4.47)$$

The effect on Bloch sphere is as shown

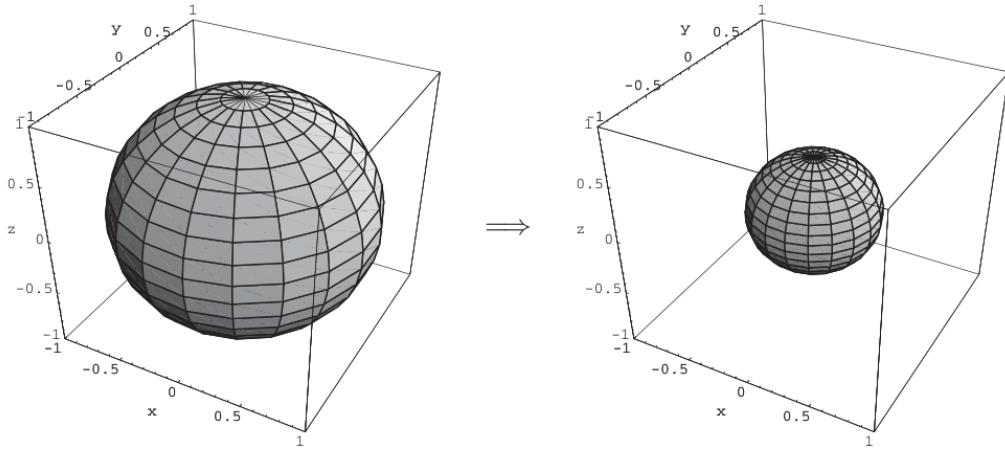


Figure 4.6: Effect of depolarizing channel on the Bloch sphere

The quantum circuit depicting the depolarizing channel is

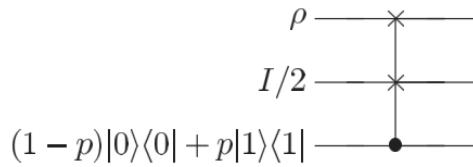


Figure 4.7: Quantum circuit representing the depolarizing channel

The third qubit is qubit is a mixture of state  $|0\rangle$  and  $|1\rangle$  with probability  $p$  and  $1 - p$  respectively. This is responsible for the second qubit getting swapped into the first qubit. In 4.47,  $I/2$  can be represented as

$$\frac{I}{2} = \frac{\rho + X\rho X + Y\rho Y + Z\rho Z}{4} \quad (4.48)$$

Thus 4.47 becomes

$$\mathcal{E}(\rho) = \left(1 - \frac{3p}{4}\right)\rho + \frac{p}{4}(X\rho X + Y\rho Y + Z\rho Z) \quad (4.49)$$

Here the operation elements will be  $\sqrt{1 - 3p/4}I$ ,  $\sqrt{p}X/2$ ,  $\sqrt{p}Y/2$ ,  $\sqrt{p}Z/2$ . Note that it is convenient to parametrize the depolarizing channel in different ways. The operation elements would differ, such as

$$\mathcal{E}(\rho) = (1 - p)\rho + \frac{p}{3}(X\rho X + Y\rho Y + Z\rho Z) \quad (4.50)$$

Here, this does nothing with probability  $1 - p$  and applies  $X$ ,  $Y$  or  $Z$  gates with probability  $p/3$  each.

The depolarizing channel can be generalized for systems with dimensionality greater than 2. Suppose a system has dimension  $d$  the quantum operation would become

$$\mathcal{E}(\rho) = \frac{pI}{d} + (1 - p)\rho \quad (4.51)$$

### 4.3.5 Amplitude damping

The process of *amplitude damping* helps us understand *energy dissipation* in quantum systems. For example, to describe dynamics of an atom emitting a photon, how a spin system at high temperature reaches equilibrium with its environment etc. What's common in these examples is described by a quantum operation known as *amplitude damping*, which can be derived as follows. Suppose we have an optical mode at state  $a|0\rangle + b|1\rangle$  which is superposition of zero or one photon. The scattering of a photon can be now described as keep a partially silvered mirror (a beamsplitter) in the path of the photon. It couples with another single optical node (its environment) according to unitary transform  $B = \exp[\theta(a^\dagger b - ab^\dagger)]$ . Where  $a, a^\dagger$  and  $b, b^\dagger$  are annihilation and creation operators for photons in two modes. The output after beamsplitter assuming there's no photon in the environment is  $B|0\rangle(a|0\rangle + b|1\rangle) = a|00\rangle + b(\cos\theta|01\rangle + \sin\theta|10\rangle)$ . Tracing out the environment gives us the quantum operation

$$\mathcal{E}_{AD}(\rho) = E_0\rho E_0^\dagger + E_1\rho E_1^\dagger \quad (4.52)$$

where  $E_k = \langle e_k | 0 | e_0 \rangle$  are

$$E_0 = \begin{bmatrix} 1 & 0 \\ 0 & \sqrt{1-\gamma} \end{bmatrix} \quad (4.53)$$

$$E_1 = \begin{bmatrix} 0 & \sqrt{\gamma} \\ 0 & 0 \end{bmatrix}, \quad (4.54)$$

the operations elements for amplitude damping.  $\gamma = \sin^2(\theta)$  can be thought of as the probability of losing a photon.  $E_1$  operation changes the state  $|1\rangle$  to  $|0\rangle$ , this corresponds to the loss of quantum of energy to the environment.  $E_0$  leaves  $|0\rangle$  unchanged but reduces the amplitude of  $|1\rangle$ , this happens because there's no loss of quantum of energy, thus the environment perceives it to be more likely that the system is in  $|0\rangle$  state.

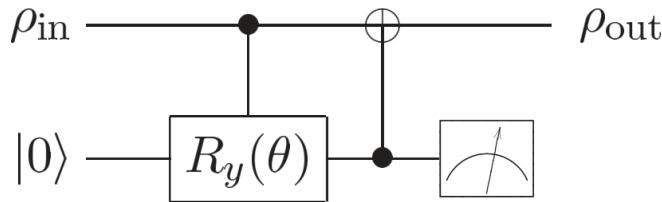


Figure 4.8: This models an amplitude damping circuit with  $\gamma = \sin^2(\theta/2)$

A general characteristic of a quantum operation is the set of states that are left invariant under the operation, as we've seen before. Here,  $|0\rangle$  is left invariant, but that is a cause of us assuming that the environment is starting at  $|0\rangle$  state, as if it were at *zero temperature*.

The quantum operation defining amplitude damping at a finite temperature is known as *generalized amplitude damping*,  $\mathcal{E}_{GAD}$  and is defined for single qubits using operation elements

$$E_0 = \sqrt{p} \begin{bmatrix} 1 & 0 \\ 0 & \sqrt{1-\gamma} \end{bmatrix} \quad E_1 = \sqrt{p} \begin{bmatrix} 0 & \sqrt{\gamma} \\ 0 & 0 \end{bmatrix} \quad (4.55)$$

$$E_2 = \sqrt{1-p} \begin{bmatrix} \sqrt{1-\gamma} & 0 \\ 0 & 1 \end{bmatrix} \quad E_3 = \sqrt{1-p} \begin{bmatrix} 0 & 0 \\ \sqrt{\gamma} & 0 \end{bmatrix} \quad (4.56)$$

and the stationary state  $\rho_\infty$  satisfies  $\mathcal{E}(\rho_\infty) = \rho_\infty$ . This GAD describes ' $T_1$ ' relaxation processes due to coupling of spin to their surrounding lattice, a large system which is in thermal equilibrium at a temperature much higher than the spin temperature. This thing is useful in NMR quantum computation.

The effect of amplitude damping on Bloch sphere is

$$(r_x, r_y, r_z) \rightarrow \left( r_x \sqrt{1 - \gamma}, r_y \sqrt{1 - \gamma}, \gamma + r_z(1 - \gamma) \right) \quad (4.57)$$

where  $\gamma$  is replaced with a time varying function like  $1 - e^{-t/T_1}$  ( $t$  is time and  $T_1$  characterizes speed of the process). Using this we can visualize the effects as *flow* on the Bloch sphere, where every point moves towards a fixed point, something like north pole where  $|0\rangle$  is located.

Similarly, generalized amplitude damping performs the transformation

$$(r_x, r_y, r_z) \rightarrow \left( r_x \sqrt{1 - \gamma}, r_y \sqrt{1 - \gamma}, \gamma(2p - 1) + r_z(1 - \gamma) \right) \quad (4.58)$$

Generalized amplitude damping is different than amplitude damping only in the location of fixed point of flow; the final state is along  $\hat{z}$  axis, at the point  $(2p - 1)$ , which is a mixed state.

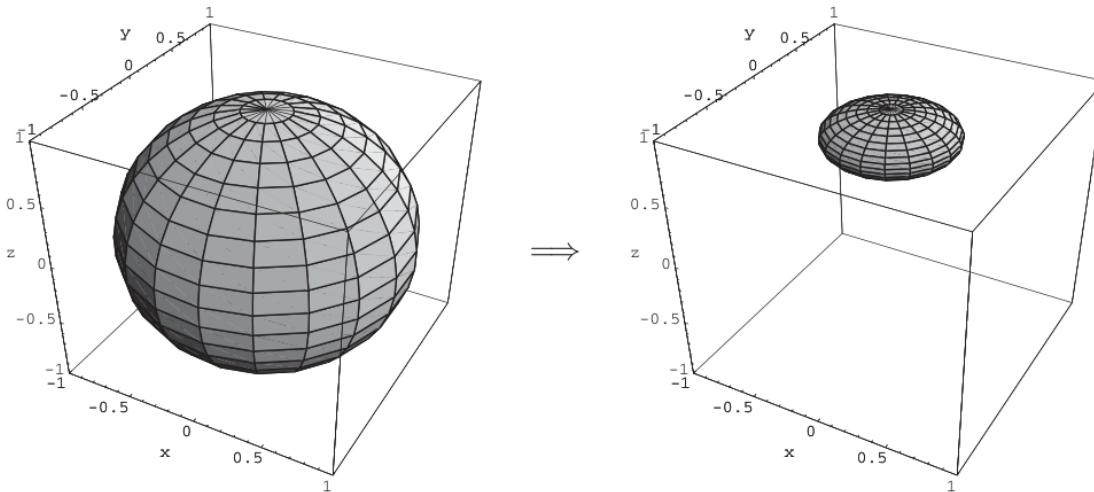


Figure 4.9: The effect of amplitude damping channel on Bloch sphere, for  $p = 0.8$ . Note how the entire sphere shrinks towards the north pole, the  $|0\rangle$  state.

### 4.3.6 Phase damping

A noise process which is uniquely quantum mechanical, describes *loss of quantum information without loss of energy*, is **phase damping**. For example it describes what happens when a photon scatters randomly as it travels through a waveguide, or how electronic states in an atom are perturbed upon interacting with distant electrical charges. The energy eigenstates don't change as a function of time but accumulate a phase proportional to the eigenvalues. When a system evolves for an amount of time not known, information about this quantum phase - the *relative* phases between the eigenstates is lost.

## **Part IV**

### **Week 4**

# Chapter 5

## Distance measures for quantum information

We'll try to quantify how two information items are similar, how information is preserved over a process etc. For this, we'll define two classes *distance measures*, *static distance* which is how close two quantum states are, and *dynamic distance* which is how well information is stored over a process. We'll define static distance properly first then build up dynamic distance with it. Two widely used distance measures currently are *trace distance* and *fidelity*.

### 5.1 Distance measures for classical information

*Hamming distance* is defined to quantify distance between two bit strings. It's equal to number of mismatching bits (eg. Hamming(1011, 0010) = 2. But this relies on labelling which we don't have in our Hilbert space arena of quantum mechanics.

We can define a classical information source as a random variable over the set of possible outcomes. So we know the probabilities of getting a given output. In this notion, we can compare two information sources with same index set using a distance measure known as *trace distance*, defined as

$$D(p_x, q_x) = \frac{1}{2} \sum_x |p_x - q_x| \quad (5.1)$$

where  $\{p_x\}$  and  $\{q_x\}$  are probability distributions of both sources. This is also known as the  *$L_1$  distance* or *Kolmogorov distance*. Trace because we use trace when we define this for quantum systems. This is a metric since it satisfies symmetry ( $D(x, y) = D(y, x)$ ) and triangle inequality ( $D(x, y) \leq D(x, z) + D(y, z)$ ).

*Fidelity*, a second measure is described by

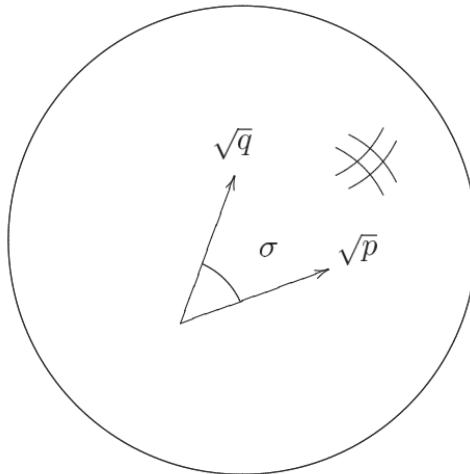
$$F(p_x, q_x) = \sum \sqrt{p_x q_x} \quad (5.2)$$

This is not a metric as shown by its physical interpretation in figure 5.1.

A physically motivated operational meaning for trace distance which can be proved is,

$$D(p_x, q_x) = \max_S |p(S) - q(S)| = \max_S \left| \sum_{x \in S} p_x - \sum_{x \in S} q_x \right| \quad (5.3)$$

which is maximum over all subsets over index set. This makes sense because  $S$  represents an event which produces maximum difference in outputs. It can also be shown that the absolute



$$F(p, q) = \sqrt{p} \cdot \sqrt{q} = \cos(\sigma)$$

Figure 5.1:  $F(p_x, q_x)$  can be seen as dot product between  $\sqrt{p}$  and  $\sqrt{q}$ , each of them lie on a unit sphere since  $\sum_x (\sqrt{p_x})^2 = 1$  and  $\sum_x (\sqrt{q_x})^2 = 1$

value is redundant, i.e

$$D(p_x, q_x) = \max_S \left( \sum_{x \in S} p_x - \sum_{x \in S} q_x \right) \quad (5.4)$$

Trace distance and fidelity are static distance measures.

For *dynamic measure* of distance, suppose a random variable  $X$  is passed through a noise channel to produce  $Y$  shown by Markov process  $X \rightarrow Y$ , a dynamic measure showing how much information is preserved is  $p(X \neq Y)$ . This can be done using trace distance, for which let's construct a copy  $\tilde{X}$  of  $X$  which is also a random variable. Now  $X$  is passed through the channel, leaving output  $Y$ . The closeness between initial perfectly correlated pair  $(\tilde{X}, X)$  and the pair  $(\tilde{X}, Y)$  can be shown to be

$$D((\tilde{X}, X), (X, Y)) = p(X \neq Y) \quad (5.5)$$

This thing we did to calculate  $p(X \neq Y)$  is unique to classical mechanics. We can't *directly* calculate quantum analogue of  $p(X \neq Y)$  if  $X$  and  $Y$  exist at different times. We use quantum entanglement to define dynamic measure similar to the construction shown above.

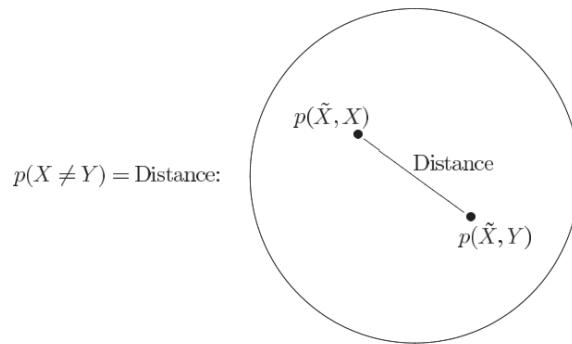


Figure 5.2: The probability of an error in the channel is equal to the trace distance between the probability distributions for  $(\tilde{X}, X)$  and  $(\tilde{X}, Y)$ .

## 5.2 How close are two quantum states

We'll see quantum generalizations of trace distance and fidelity.

### 5.2.1 Trace distance

*Trace distance* between two quantum states  $\rho$  and  $\sigma$  is defined as

$$D(\rho, \sigma) = \frac{1}{2} \text{tr}|\rho - \sigma| \quad (5.6)$$

here  $|A| = \sqrt{A^\dagger A}$ . If  $\rho$  and  $\sigma$  commute (matrix multiplication), this reduces to classical trace distance. Suppose  $\rho$  and  $\sigma$  commute, then

$$\rho = \sum_i r_i |i\rangle\langle i|; \quad \sigma = \sum_i s_i |i\rangle\langle i| \quad (5.7)$$

for orthogonal basis  $|i\rangle$ . Then trace distance becomes

$$D(\rho, \sigma) = \frac{1}{2} \text{tr}|\rho - \sigma| \quad (5.8)$$

$$= \frac{1}{2} \text{tr}|(r_i - s_i)| \quad (5.9)$$

$$= D(r_i, s_i) \quad (5.10)$$

For qubits in Bloch sphere representation, it becomes nicer. Let

$$\rho = \frac{1 + \vec{r}\vec{\sigma}}{2}; \quad \sigma = \frac{1 + \vec{s}\vec{\sigma}}{2} \quad (5.11)$$

then trace distance between  $\rho$  and  $\sigma$  is

$$D(\rho, \sigma) = \frac{|\vec{r} - \vec{s}|}{2} \quad (5.12)$$

using the fact that  $(\vec{r} - \vec{s}) \cdot \vec{\sigma}$  has eigen values  $\pm |\vec{r} - \vec{s}|$ . It can be shown that

$$D(\rho, \sigma) = \max_P \text{tr}(P(\rho - \sigma)) \quad (5.13)$$

where maximum is over all positive operators  $P$ ,  $P \leq I$ .

**Theorem 11.** Let  $\{E_m\}$  me a POVM, with  $p_m \equiv \text{tr}(\rho E_m)$  and  $q_m \equiv \text{tr}(\sigma E_m)$  be the probability of getting output labelled by  $m$ . Then

$$D(\rho, \sigma) = \max_{\{E_m\}} D(p_m, q_m) \quad (5.14)$$

where maximization is over all POVMs  $\{E_m\}$ .

It can be shown that this trace distance is also a metric. Here goes another nice theorem

**Theorem 12.** Suppose  $\mathcal{E}$  is a trace preserving quantum operation. Let  $\rho, \sigma$  be density operators. Then

$$D(\mathcal{E}(\rho), \mathcal{E}(\sigma)) \leq D(\rho, \sigma) \quad (5.15)$$

Thus there exists no physical processes which increases distance between two states.

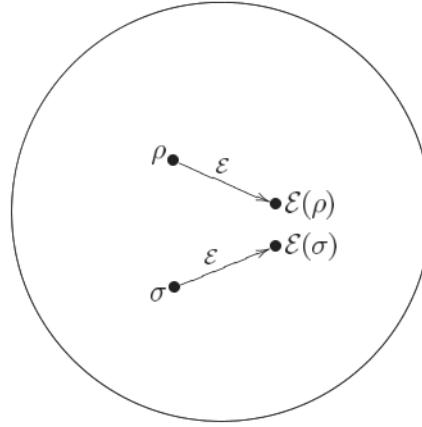


Figure 5.3: Trace-preserving quantum operations cause a contraction on the space of density operators.

It can also be showed that

$$D(\rho^A, \sigma^A) \leq D(\rho^{AB}, \sigma^{AB}) \quad (5.16)$$

intuitively, it's like things are less differentiable if corresponding parts of them are covered.

**Theorem 13 (Strong convexity of trace distance).** Let  $\{p_i\}$  and  $\{q_i\}$  be two probability distributions over the same index set, and  $\rho_i, \sigma_i$  be density operators, also with indices from same index set, then

$$D\left(\sum_i p_i \rho_i, \sum_i q_i \sigma_i\right) \leq D(p_i, q_i) + \sum_i p_i D(\rho_i, \sigma_i) \quad (5.17)$$

where  $D(p_i, q_i)$  is the trace distance between the probability distributions  $p_i$  and  $q_i$ .

using this we can show convexity properties of trace distance,

$$D\left(\sum_i p_i \rho_i, \sigma\right) \leq \sum_i p_i D(\rho_i, \sigma) \quad (5.18)$$

Also, any trace-preserving quantum operation  $\mathcal{E}$  has a fixed point  $\rho$  such that  $\mathcal{E}(\rho) = \rho$ .

### 5.2.2 Fidelity

*Fidelity* is not a metric on density operators but does give rise to a useful metric. Fidelity of state  $\rho$  and  $\sigma$  is defined as

$$F(\rho, \sigma) \equiv \text{tr} \sqrt{\rho^{1/2} \sigma \rho^{1/2}} \quad (5.19)$$

when  $\rho$  and  $\sigma$  commute, then

$$\rho = \sum_i p_i |i\rangle \langle i|; \quad \sigma = \sum_i s_i |i\rangle \langle i| \quad (5.20)$$

for orthonormal basis  $|i\rangle$ , we see that

$$F(\rho, \sigma) = \text{tr} \sqrt{\sum_i r_i s_i |i\rangle \langle i|} \quad (5.21)$$

$$= \text{tr} \left( \sum_i \sqrt{r_i s_i} |i\rangle \langle i| \right) \quad (5.22)$$

$$= \sum_i \sqrt{r_i s_i} \quad (5.23)$$

$$= F(r_i, s_i) \quad (5.24)$$

thus, when  $\rho$  and  $\sigma$  commute, quantum fidelity reduces to classical fidelity between eigenvalue distributions  $r_i, s_i$  of  $\rho, \sigma$  respectively. Also the fidelity between a pure state  $|\psi\rangle$  and an arbitrary state  $\rho$  is

$$F(|\psi\rangle, \rho) = \text{tr} \sqrt{\langle \psi | \rho | \psi \rangle} \quad (5.25)$$

$$= \sqrt{\langle \psi | \rho | \psi \rangle} \quad (5.26)$$

i.e fidelity is square root of overlap between  $|\psi\rangle$  and  $\rho$ .

There's no similar Bloch sphere representation for fidelity but it does satisfy similar properties like *invariance under unitary transformation* i.e

$$F(U\rho U^\dagger, U\sigma U^\dagger) = F(\rho, \sigma) \quad (5.27)$$

**Theorem 14 (Uhlmann's theorem).** Suppose  $\rho, \sigma$  be states of a system  $Q$ , Let  $R$  be a copy of  $Q$ . Then

$$F(\rho, \sigma) = \max_{|\psi\rangle, |\varphi\rangle} |\langle \psi | \varphi \rangle| \quad (5.28)$$

where the maximization is over all purifications  $|\psi\rangle$  of  $\rho$  and  $|\varphi\rangle$  of  $\sigma$ .

**Lemma 1.** Let  $A$  be any operator,  $U$  be unitary. Then

$$|\text{tr}(AU)| \leq |\text{tr}(A)| \quad (5.29)$$

equality is when  $U = V^\dagger$  when  $A = |A|V$  is polar decomposition of  $A$ .

By Uhlmann's formula it can be shown that fidelity is *symmetric* in its inputs i.e  $F(\rho, \sigma) = F(\sigma, \rho)$  and also that it's bound,  $0 \leq F(\rho, \sigma) \leq 1$ . Also  $\rho = \sigma \implies F(\rho, \sigma) = 1$  else it's less than 1. Quantum fidelity is related to classical fidelity in the following way:

$$F(\rho, \sigma) = \min_{E_m} F(p_m, q_m) \quad (5.30)$$

where  $E_m$  are POVMs,  $p_m = \text{tr}(\rho E_m)$ ,  $q_m = \text{tr}(\sigma E_m)$  are probability distributions corresponding to the POVM  $\{E_m\}$ .

We can define *angle* between states  $\rho$  and  $\sigma$  by

$$A(\rho, \sigma) = \arccos F(\rho, \sigma) \quad (5.31)$$

This is non-negative, symmetric in inputs and zero when  $\rho = \sigma$ . It can be shown to satisfy triangle inequality, and thus is a metric.

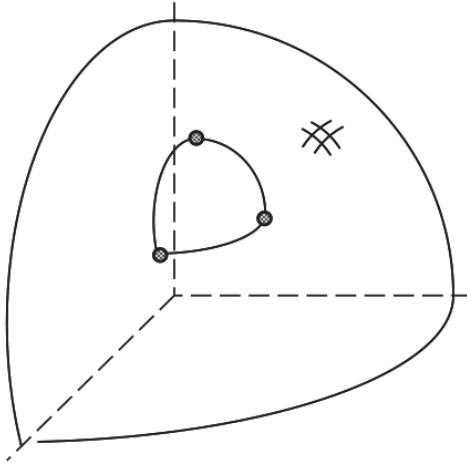


Figure 5.4: Angle between unit vectors is a metric

Qualitatively, fidelity behaves like “upside down” of trace distance. It decreases as two states become more distinguishable and vice versa. Fidelity is *non decreasing* as shown by the theorem

**Theorem 15 (Monotonicity of fidelity).** Suppose  $\rho, \sigma$  be density operators, and  $\mathcal{E}$  is a trace-preserving operation, then

$$F(\mathcal{E}(\rho), \mathcal{E}(\sigma)) \geq F(\rho, \sigma) \quad (5.32)$$

Angles we've defined above follow **contractivity** which is  $A(\mathcal{E}(\rho), \mathcal{E}(\sigma)) \leq A(\rho, \sigma)$ , there's also **strong concavity** of fidelity like trace distance

**Theorem 16 (Strong concavity of fidelity).** Let  $p_i, q_i$  be probability distributions over the same index set,  $\rho_i, \sigma_i$  also indexed by the same index set. Then

$$F\left(\sum_i p_i \rho_i, q_i \sigma_i\right) \geq \sum_i \sqrt{p_i q_i} F(\rho_i, \sigma_i) \quad (5.33)$$

### 5.2.3 Relationships between distance measures

Despite their different forms, trace distance and fidelity are closely related. In pure states, they're completely equivalent to each other, to see this consider two pure state  $|a\rangle$  and  $|b\rangle$ , using Gram-schmidt we can find states  $|0\rangle$  and  $|1\rangle$  such that  $|a\rangle = |0\rangle$  and  $|b\rangle = \cos \theta |0\rangle + \sin \theta |1\rangle$ . Then  $F(|a\rangle, |b\rangle) = |\cos \theta|$  and

$$D(|a\rangle, |b\rangle) = \frac{1}{2} \text{tr} \begin{vmatrix} 1 - \cos^2 \theta & -\cos \theta \sin \theta \\ -\cos \theta \sin \theta & -\sin^2 \theta \end{vmatrix} \quad (5.34)$$

$$= |\sin \theta| \quad (5.35)$$

$$= \sqrt{1 - F(|a\rangle, |b\rangle)^2} \quad (5.36)$$

For mixed states  $\rho, \sigma$  let  $|\psi\rangle$  and  $|\varphi\rangle$  be purifications of  $\rho, \sigma$  such that  $F(\rho, \sigma) = |\langle \psi | \varphi \rangle| = F(|\psi\rangle, |\varphi\rangle)$ . Since trace distance is non-increasing under the partial trace,

$$D(\rho, \sigma) \leq D(|\psi\rangle, |\varphi\rangle) \quad (5.37)$$

$$= \sqrt{1 - F(\rho, \sigma)^2} \quad (5.38)$$

It can be further shown that

$$1 - F(\rho, \sigma) \leq D(\rho, \sigma) \leq \sqrt{1 - F(\rho, \sigma)^2} \quad (5.39)$$

### 5.3 How well does a quantum channel preserve information?

We'll try to quantify how much that state has "changed" from  $|\psi\rangle$  to  $\mathcal{E}(|\psi\rangle\langle\psi|)$  due to an operation  $\mathcal{E}$  using the distance measures discussed. Let's look at a simple example how a state changes under depolarizing channel using fidelity

$$F(|\psi\rangle, \mathcal{E}(|\psi\rangle\langle\psi|)) = \sqrt{\langle\psi| \left( p\frac{I}{2} + (1-p)|\psi\rangle\langle\psi| \right) |\psi\rangle} \quad (5.40)$$

$$= \sqrt{1 - \frac{p}{2}} \quad (5.41)$$

it can be clearly seen when we reduce  $p$ , fidelity get's close to one hence the states becoming more indistinguishable. There's nothing special about fidelity that trace distance doesn't have but we'll stick to this. But in real quantum systems, we don't even know initial state  $|\psi\rangle$  in advance so we try to *minimize* fidelity to get worst-case measure

$$F_{\min}(\mathcal{E}) \equiv \min_{|\psi\rangle} F(|\psi\rangle, \mathcal{E}(|\psi\rangle\langle\psi|)) \quad (5.42)$$

for depolarizing channel it doesn't matter and it's just  $\sqrt{1-p/2}$ . But for phase damping channel

$$\mathcal{E}(\rho) = p\rho + (1-p)Z\rho Z \quad (5.43)$$

The fidelity will be

$$F(|\psi\rangle, \mathcal{E}(|\psi\rangle\langle\psi|)) = \sqrt{\langle\psi| (p|\psi\rangle\langle\psi| + (1-p)Z|\psi\rangle\langle\psi|Z) |\psi\rangle} \quad (5.44)$$

$$= \sqrt{p + (1-p)\langle\psi|Z|\psi\rangle^2} \quad (5.45)$$

Now this minimizes when  $\frac{|0\rangle+|1\rangle}{\sqrt{2}}$ . Thus for phase damping, minimum fidelity becomes

$$F_{\min}(\mathcal{E}) = \sqrt{p} \quad (5.46)$$

We haven't considered mixed states, but using joint concavity of fidelity it can be shown that allowing mixed states doesn't change  $F_{\min}$ . Now, we're not only interesting in saving quantum states while transmitting them but we should also look up on when we're computing for example implementing a quantum gate described by  $U$ . Thus we define *gate-fidelity* as

$$F(U, \mathcal{E}) = \min_{|\psi\rangle} F(U|\psi\rangle, \mathcal{E}(|\psi\rangle\langle\psi|)) \quad (5.47)$$

For example, if we're implementing NOT gate, but instead implement a noisy operation  $\mathcal{E}(\rho) = (1-p)X\rho X + pZ\rho Z$ , then gate fidelity is

$$F(X, \mathcal{E}) = \min_{|\psi\rangle} \sqrt{\langle\psi| X ((1-p)X|\psi\rangle\langle\psi|X + pZ|\psi\rangle\langle\psi|Z) X |\psi\rangle} \quad (5.48)$$

$$= \min_{|\psi\rangle} = \sqrt{(1-p) + p\langle\psi|Y|\psi\rangle^2} \quad (5.49)$$

$$= \sqrt{1-p}. \quad (5.50)$$

### 5.3.1 Quantum sources of information and the entanglement fidelity

Even though we've talked a lot about quantum sources we haven't actually defined it. One interesting definition is imagining a stream of quantum systems being produced by some physical process, with their states represented as  $\rho_{X_1}, \rho_{X_2}, \dots$ , where  $X_j$  are identically distributed random variables and  $\rho_j$  is some fixed set of operators. For example, a stream of qubits preparing  $|0\rangle$  with 1/2 probability and  $(|0\rangle + |1\rangle)/\sqrt{2}$  with another 1/2 of probability.

This *ensemble* notion of quantum source leads to notion of *ensemble average fidelity* which captures the idea that source is well preserved under the action of a noisy channel described by trace preserving operator  $\mathcal{E}$ , as

$$F = \sum_j p_j F(\rho_j, \mathcal{E}(\rho_j))^2, \quad (5.51)$$

Here,  $F = 1$  if and only if  $\mathcal{E}(\rho_j) = \rho_j$  for all  $j$  such that  $p_j > 0$

The second notion of quantum source is motivated by the fact that a channel preserving information also preserves *entanglement* well. We'll use quantum analogue of what we've considered in 5.5. We'll suppose a quantum system  $Q$  entangled with a fictitious system  $R$ , such that the joint state  $RQ$  is pure. It turns out the result doesn't depend on how this purification is done. Then the system  $Q$  is subjected to dynamics described by  $\mathcal{E}$ .

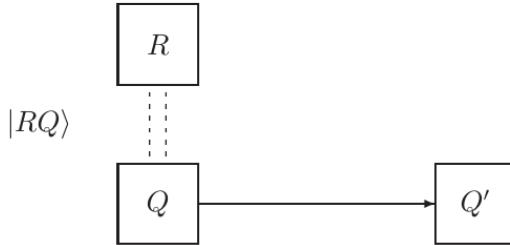


Figure 5.5: The  $RQ$  picture of quantum channel. The initial state of  $RQ$  is a pure state.

How well is the entanglement preserved by the quantum operation  $\mathcal{E}$  is described by *entanglement fidelity*  $F(\rho, \mathcal{E})$  which is defined for trace preserving operations  $\mathcal{E}$  by

$$F(\rho, \mathcal{E}) \equiv F(RQ, R'Q')^2 \quad (5.52)$$

$$= \langle RQ | [(\mathcal{I}_R \otimes \mathcal{E})(|RQ\rangle \langle RQ|)] | RQ \rangle \quad (5.53)$$

where prime implies after the operation. Entanglement fidelity close to 1 implies well preserved state and vice versa. One of the important properties of entanglement fidelity is that there's a nice formula to calculate it exactly. Suppose  $E_i$  is a set of operation elements for  $\mathcal{E}$ . Then

$$F(\rho, \mathcal{E}) = \langle RQ | \rho^{R'Q'} | RQ \rangle = \sum_i |\langle RQ | E_i | RQ \rangle|^2. \quad (5.54)$$

Suppose  $|RQ\rangle = \sum_j \sqrt{p_j} |j\rangle |j\rangle$ , where  $\rho = \sum_j |j\rangle \langle j|$ . Then it can be proved that

$$\langle RQ | E_i | RQ \rangle = \text{tr}(E_i \rho) \quad (5.55)$$

Substituting this in equation 5.54 we get a useful formula

$$F(\rho, \mathcal{E}) = \sum_i |\text{tr}(\rho E_i)|^2 \quad (5.56)$$

Surprisingly, these two notions are closely related too. This can be seen in two useful properties. First, the entanglement fidelity is a lower bound on the square of the static fidelity between the output and input to the process,

$$F(\rho, \mathcal{E}) \leq [F(\rho, \mathcal{E}(\rho))]^2 \quad (5.57)$$

This states that it's harder to preserve entanglement than it is compared to pure states. The second property is that it is a *convex* function of  $\rho$ . i.e

$$F\left(\sum_j p_j \rho_j, \mathcal{E}\right) \leq \bar{F} \quad (5.58)$$

This shows that any channel  $\mathcal{E}$  which does a good job of preserving entanglement between a source system  $\rho$  and reference system will automatically preserve the ensemble source described by probabilities  $p_j$  and states  $\rho_j$  such that  $\sum_j p_j \rho_j = \rho$ .

I'll try to conclude with most important points yet.

1.  $0 \leq F(\rho, \mathcal{E}) \leq 1$ . From properties of static fidelity.
2. Entanglement fidelity is linear in quantum operation input. This is immediate too.
3. For pure state inputs, entanglement fidelity is square of static fidelity between input and output,

$$F(|\psi\rangle, \mathcal{E}) = F(|\psi\rangle, \mathcal{E}(|\psi\rangle\langle\psi|))^2. \quad (5.59)$$

4.  $F(\rho, \mathcal{E}) = 1$  if and only if for all pure state  $|\psi\rangle$  in support of  $\rho$

$$\mathcal{E}(|\psi\rangle\langle\psi|) = |\psi\rangle\langle\psi|. \quad (5.60)$$

5. Suppose that  $\langle\psi|\mathcal{E}(|\psi\rangle\langle\psi|)|\psi\rangle \leq 1 - \eta$  for all  $|\psi\rangle$  in support of  $\rho$ , for some  $\eta$  then  $F(\rho, \mathcal{E}) \geq 1 - (3\eta/2)$ .

**Part V**

**Week 5**

# Chapter 6

## Entropy and information

We know that *entropy* measures uncertainty of state of a physical system. We'll look into it's definitions and properties deeply both in classical and quantum information theory.

### 6.1 Shannon entropy

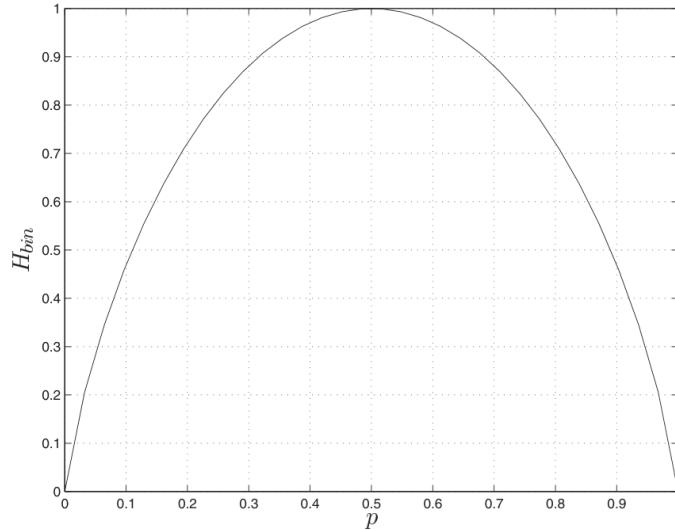
It's the key concept of classical information theory. It can be viewed in two complementary views, first, suppose we have a random variable  $X$ , how much information we *gain* on average if we know the value of  $X$ . Second, it's the measure of *uncertainty* before we know the value of  $X$ .

Information content of a random variable doesn't depend on labels. For example, a random variable giving 'heads' and 'tails' with probability  $1/4$  and  $3/4$  contains same information as a random variable giving '0' and '1' with probability  $1/4$  and  $3/4$  respectively. Thus, entropy also doesn't depend on labels. *Shannon entropy* associated with a probability distribution  $p_1, \dots, p_n$  is given by

$$H(X) \equiv H(p_1, \dots, p_n) \equiv - \sum_x p_x \log p_x \quad (6.1)$$

Note that by  $\log$  we mean base 2, by  $\ln$  we mean base  $e$ . Since a never occurring event with probability  $p_i = 0$  should never affect entropy, thus we agree that  $0 \log 0 = 0$  even though we mean  $\lim_{x \rightarrow 0} x \log x = 0$ . There's a nice reason why entropy is defined this way. We have to *quantify the resources needed to store information*. Let's suppose a source producing symbols  $X_1, X_2, \dots$  of independent *identically* distributed random variables, minimal amount of resources required to store the information produced by source such that it can be reconstructed turns out to be just entropy i.e  $H(X) \equiv H(X_1) \equiv H(X_2) = \dots$ , this result is known as *Shannon's noiseless coding theorem*.

I suggest looking at an example to understand this. Suppose our source produces 1, 2, 3, 4 which are symbols but with probability  $1/2, 1/4, 1/8$  and  $1/8$  respectively. We don't need two bits to store output of each use of source. We can compress this usage, by representing 1 as 0, 2 as 10, 3 as 110 and 4 as 111, with this we can reconstruct the message, for eg. 11010111 as 324 too. We need on an average  $\frac{1}{2} \cdot 1 + \frac{1}{4} \cdot 2 + \frac{1}{8} \cdot 3 + \frac{1}{8} \cdot 3 = 7/4$  bits only! It also turns out that this is just the entropy of the source  $H(X) = -1/2 \log(1/2) - 1/4 \log(1/4) - 1/8 \log(1/8) - 1/8 \log(1/8) = 7/4$ . It's also true that average information gained by an event out of a probability distribution  $p_1, \dots, p_n$  is  $k \sum_i p_i \log p_i$  for some constant  $k$ .

Figure 6.1: Binary entropy function  $H(p)$ 

## 6.2 Basic properties of entropy

### 6.2.1 The binary entropy

The entropy of a two-outcome random variable is so special that we name it as *binary entropy*

$$H_{\text{bin}}(X) \equiv -p \log p - (1-p) \log(1-p) \quad (6.2)$$

where  $p$  and  $1-p$  are probabilities of outcome. Its graph is shown in figure 6.1

A really good intuition comes when we consider this example of Alice having two biased coins one from US and another from Australia. Suppose the US coin gives head with probability  $p_U$  and the Australian coin with probability  $p_A$  and Alice flips one of those coins with probability  $q$  for US coin and  $1-q$  for Australian coin and she tells final ‘head’ or ‘tail’ to Bob. The information Bob gets is atleast the average information he gets from flipping one coin. Hence

$$H(qp_U + (1-q)p_A) \geq qH(p_U) + (1-q)H(p_A) \quad (6.3)$$

It's greater because Bob gets *more* information about the country sometimes. For eg. if  $p_U = 0.1$  and  $p_A = 0.9$  and Alice says head, Bob would understand that it's more probable to be an aussie coin. This leads us into a very important property of entropy, *concavity*. Which can be seen in binary case in the figure 6.1.

### 6.2.2 The relative entropy

The *relative entropy* is useful entropy-like measure describing closeness of two probability distributions  $p(x)$  and  $q(x)$  over the same index set  $x$ , defined by

$$H(p(x) \parallel q(x)) \equiv \sum_x p(x) \log \frac{p(x)}{q(x)} \equiv -H(X) - \sum_x p(x) \log q(x) \quad (6.4)$$

The following theorem motivates us to understand it as a distance measure

**Theorem 17 (Non-negativity of relative entropy).** *The relative entropy is non-negative,  $H(p(x) \parallel q(x)) \geq 0$  with equality if and only if  $p(x) = q(x)$  for all  $x$ .*

The proof is simple using the definition of relative entropy and  $\log x \ln 2 \leq x - 1$  with equality only when  $x = 1$ . Relative entropy itself is not very useful but used to study entropy, for example let  $q(x)$  be uniform probability distribution over index set  $X$  of size  $d$  then for a probability distribution  $p(x)$  over  $X$ , we have

$$H(p(x) \parallel q(x)) = H(p(x) \parallel 1/d) = -H(X) - \sum_x p(x) \log(1/d) = \log d - H(X) \quad (6.5)$$

and hence by theorem 17 we have  $\log d - H(X) \geq 0$  thus  $H(X) \leq \log d$ , a nice property hence a theorem!

**Theorem 18.** *Suppose  $X$  is a random variable with  $d$  outcomes. Then  $H(X) \leq \log d$ , with equality if and only if  $X$  is uniformly distributed over those  $d$  outcomes.*

### 6.2.3 Conditional entropy and mutual information

For two random variables  $X, Y$  we'll try to understand how information of  $X$  is related to that of  $Y$ . First we define *joint entropy* of  $X$  and  $Y$  as

$$H(X, Y) \equiv - \sum_{x,y} p(x, y) \log p(x, y) \quad (6.6)$$

The *entropy of  $X$  conditional on knowing  $Y$*  is similar to that in probability, it's the uncertainty that we have in the value of  $X$  given that we know the value of  $Y$ . It's defined by

$$H(X|Y) \equiv H(X, Y) - H(Y) \quad (6.7)$$

A second quantity *mutual information content of  $X$  and  $Y$*  measures how much information  $X$  and  $Y$  have in common. It's very similar to intersection, given we take  $H(X, Y)$  as information content of both  $X$  and  $Y$  which is counted twice. Thus, mutual information content is

$$H(X : Y) = H(X) + H(Y) - H(X, Y) \quad (6.8)$$

The relation  $H(X : Y) = H(X) - H(X|Y)$  also comes in handy. Let's see some more simple relationships between different entropies.

#### Theorem 1 (Basic properties of Shannon entropy)

1.  $H(X, Y) = H(Y, X), H(X : Y) = H(Y : X)$
2.  $H(Y|X) \geq 0$  and thus  $H(X : Y) \leq H(Y)$  with equality if and only if  $Y$  is a function of  $X$ ,  $Y = f(X)$ .
3.  $H(X) \leq H(X, Y)$  with equality if and only if  $Y$  is a function of  $X$ .
4. **Subadditivity:**  $H(X, Y) \leq H(X) + H(Y)$  with equality if and only if  $X$  and  $Y$  are independent variables.
5.  $H(Y|X) \leq H(Y)$  and thus  $H(X : Y) \geq 0$  with equality in each if and only if  $X$  and  $Y$  are independent distributions.

6. **Strong subadditivity:**  $H(X, Y, Z) + H(Y) \leq H(X, Y) + H(Y, Z)$ , with equality if and only if  $Z \rightarrow Y \rightarrow X$  forms a markov chain.

7. **Conditioning reduces entropy:**  $H(X|Y, Z) \leq H(X|Y)$ .

All this can be easily understood by a Venn diagram for entropy.

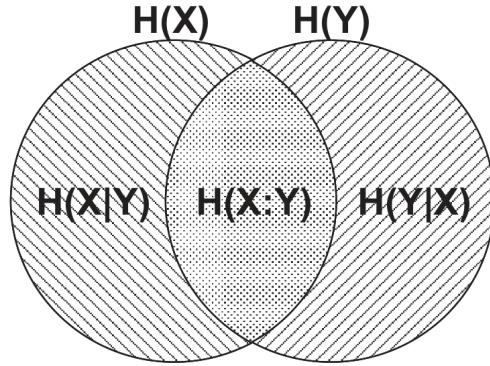


Figure 6.2: Relationships between different entropies

Let's finally consider a simple and useful chaining rule for conditional entropies.

**Theorem 19.** Let  $X_1, \dots, X_n$  and  $Y$  be any set of random variables. Then

$$H(X_1, \dots, X_n|Y) = \sum_{i=1}^n H(X_i|Y, X_1, \dots, X_{i-1}) \quad (6.9)$$

*Proof.* Use induction with base case  $n = 2$ . □

#### 6.2.4 The data processing inequality

This states that information from the output of a source can only *decrease with time*; once information is lost, it's gone forever. The intuitive notion of *information processing* is captured by *Markov chains* of random variables. A Markov chain is a sequence  $X_1 \rightarrow X_2 \rightarrow \dots$  such that  $X_{n+1}$  is independent of  $X_1, X_2, \dots, X_{n-1}$  if  $X_n$  is known. In other words

$$p(X_{n+1} = x_{n+1}|X_n = x_n, \dots, X_1 = x_1) = p(X_{n+1} = x_{n+1}|X_n = x_n) \quad (6.10)$$

**Theorem 20.** Suppose  $X \rightarrow Y \rightarrow Z$  is a Markov chain. Then

$$H(X) \geq H(X : Y) \geq H(X : Z). \quad (6.11)$$

Moreover, this first inequality becomes equal if and only if, given  $Y$ , it's possible to reconstruct  $X$ .

This theorem is intuitive, as it says that if there's some noise introduced in  $X$  (which leads to Markov chain  $X \rightarrow Y$ ) then the mutual information between  $X$  and  $Y$  (which is what we want) is irretrievably lost, since it's not possible to reconstruct  $X$  from  $Y$  (noise).

If  $X \rightarrow Y \rightarrow Z$  is a Markov chain, then so is  $Z \rightarrow Y \rightarrow X^1$ . Thus with above conditions, we see that

$$H(Z : Y) \geq H(Z : X) \quad (6.12)$$

This thing is also known as *data pipelining inequality* i.e any information  $Z$  shares with  $X$  is also shared by  $Y$ . The information is *pipelined* from  $X$  to  $Y$  to  $Z$ .

<sup>1</sup>Can be proved by using  $P(X|Y) = P(X, Y)/P(Y)$

## 6.3 Von Neumann entropy

Shannon entropy is for classical probability distributions, Von Neumann defined this notion of *entropy* for quantum state with density operator  $\rho$  as

$$S(\rho) \equiv -\text{tr}(\rho \log \rho) \quad (6.13)$$

where  $\log$  still means base 2 and  $0 \log 0 \equiv 0$ . This can be simplified if we know eigenvalues  $\lambda_x$  of  $\rho$  as

$$S(\rho) \equiv -\sum_x \lambda_x \log \lambda_x \quad (6.14)$$

This is very useful for calculations. We will refer *entropy* to either Shannon or Von Neumann entropy based on context.

### 6.3.1 Quantum relative entropy

Similar to Shannon entropy, *relative entropy* between two state  $\rho$  and  $\sigma$  is defined by

$$S(\rho \| \sigma) \equiv \text{tr}(\rho \log \rho) - \text{tr}(\rho \log \sigma) \quad (6.15)$$

as in classical case, this can sometimes shoot upto infinity. Particularly, relative entropy is defined to be  $+\infty$  if the kernel of  $\sigma$  (vector space spanned by eigenvectors of  $\sigma$  with eigenvalue 0) has non-trivial intersection with the support of  $\rho$  (vector space spanned by eigenvectors of  $\rho$  with non-zero eigenvalues). A familiar inequality as a theorem is given.

**Theorem 21 (Klein's inequality).** *The quantum relative entropy is non-negative. i.e*

$$S(\rho \| \sigma) \geq 0, \quad (6.16)$$

*with equality if and only if  $\rho = \sigma$ .*

### 6.3.2 Basic properties of entropy

These are few interesting properties of von Neumann entropy:

**Theorem 2 (Basic properties of von Neumann entropy)**

1. The entropy is non-negative. The entropy is zero if and only if the state is pure.
2. In a  $d$ -dimensional Hilbert space, entropy is atmost  $\log d$ . It's equal if and only if the system is completely mixed state  $I/d$ .
3. If a composite system  $AB$  is pure, then  $S(A) = S(B)$ .
4. Suppose  $p_i$  are probabilities and  $\rho_i$  have support on orthogonal subspaces. Then

$$S\left(\sum_i p_i \rho_i\right) = H(p_i) + \sum_i p_i S(\rho_i) \quad (6.17)$$

5. **Joint entropy theorem:** Suppose  $p_i$  are probabilities and  $|i\rangle$  are orthogonal states for a system  $A$ , and  $\rho_i$  is any set of density operators for another system,  $B$ . Then

$$S\left(\sum_i p_i |i\rangle \langle i| \otimes \rho_i\right) = H(p_i) + \sum_i p_i S(\rho_i). \quad (6.18)$$

Again similarly, the *joint entropy*  $S(A, B)$  for a composite system of two components  $A$  and  $B$  is defined obviously as  $S(A, B) \equiv -\text{tr}(\rho^{AB} \log \rho^{AB})$ , we also define conditional entropy and mutual information as

$$S(A|B) = S(A, B) - S(B) \quad (6.19)$$

$$S(A : B) = S(A) + S(B) - S(A, B) \quad (6.20)$$

$$= S(A) - S(A|B) = S(B) - S(B|A). \quad (6.21)$$

But there are some properties which aren't carried over from Shannon to von Neumann entropy. One such is  $H(X, Y) \geq H(X)$  which is intuitive since we have more information if we know about  $X$  and  $Y$  compared to if we only know about  $X$ . But this isn't true for von Neumann entropy, consider a system  $AB$  of two qubits in entangled state  $(|00\rangle + |11\rangle)/\sqrt{2}$ . Since this is pure,  $S(A, B) = 0$ . But its component  $A$  is  $I/2$ , and has entropy 1 defying what we've seen above. Another way of seeing this is  $S(B|A) = S(A, B) - S(A)$  is negative.

### 6.3.3 Measurements and entropy

We'll try to understand how measurements change the entropy of a system, if they do. This change of course depends on the type of measurement we make. Suppose we do a projective measurement on system  $\rho$  described by projectors  $P_i$ , the state of the system is

$$\rho' = \sum_i P_i \rho P_i \quad (6.22)$$

The following theorem shows the entropy never decreases

**Theorem 22 (Projective measurements increase entropy).** *Suppose  $P_i$  is a complete set of orthogonal projectors and  $\rho$  is a density operator. Then the entropy of the state  $\rho' \equiv \sum_i P_i \rho P_i$  of the system after measurement is atleast as great as the original entropy,*

$$S(\rho') \geq S(\rho) \quad (6.23)$$

with equality if and only if  $\rho' = \rho$ .

### 6.3.4 Subadditivity

If two systems  $A$  and  $B$  have a joint state  $\rho^{AB}$ . Then the joint entropy for the two systems satisfy

$$S(A, B) \leq S(A) + S(B) \quad (6.24)$$

$$S(A, B) \geq |S(A) - S(B)| \quad (6.25)$$

The first relation is known as *subadditivity* inequality<sup>2</sup> for von Neumann entropy, with equality if and only if  $A$  and  $B$  are correlated, i.e  $\rho^{AB} = \rho^A \otimes \rho^B$ . Second one is known as *triangle* inequality or *Araki-Lieb* inequality. It's the quantum analogue of  $H(X, Y) \geq H(X)$ .

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<sup>2</sup>This can be proved simply using Klein's inequality.

### 6.3.5 Concavity of the entropy

The entropy is a *concave* function in its inputs. That is, given the probabilities  $p_i$ , and corresponding density operators  $\rho_i$ , entropy satisfies the inequality<sup>3</sup>

$$S\left(\sum_i p_i \rho_i\right) \geq \sum_i p_i S(\rho_i). \quad (6.26)$$

A good intuition is that  $\sum_i p_i \rho_i$  expresses a quantum state which is in an uncertain state  $\rho_i$  with probability  $p_i$ , this uncertainty of index  $i$  of state adds to the average uncertainty we have of each state  $\rho_i$  which is  $\sum_i p_i S(\rho_i)$ . Equality holds if and only if all the states  $\rho_i$  for which  $p_i > 0$  are identical; i.e, entropy is strictly concave in its inputs!

### 6.3.6 The entropy of a mixture of quantum states

The flip side of concavity is that it provides an upper bound of entropy of a mixture of quantum states too. It's shown by the following theorem.

**Theorem 23.** Suppose  $\rho = \sum_i p_i \rho_i$ , where  $p_i$  are some set of probabilities, and  $\rho_i$  are density operators. Then

$$S(\rho) \leq \sum_i p_i S(\rho_i) + H(p_i), \quad (6.27)$$

with equality if and only if the states  $\rho_i$  have support on orthogonal subspaces.

This along with concavity completes our bound on entropy of a mixture of quantum states,

$$\sum_i p_i S(\rho_i) \leq S\left(\sum_i p_i \rho_i\right) \leq \sum_i p_i S(\rho_i) + H(p_i). \quad (6.28)$$

The intuition for this comes from the fact that we know too much information if we know on average about  $\rho_i$  and  $p_i$  also.

## 6.4 Strong subadditivity

This is an extension of subadditivity and triangle inequalities to three systems. The inequality states that for three quantum systems  $A, B, C$ ,

$$S(A, B, C) + S(B) \leq S(A, B) + S(B, C) \quad (6.29)$$

### 6.4.1 Proving strong subadditivity

The proof is based on an important result, the *Lieb's theorem*. For which, let's define few things. A function  $f(A, B)$  which is a real-valued function of two matrices, is said to be *jointly concave* in  $A, B$  if for all  $0 \leq \lambda \leq 1$ ,

$$f(\lambda A_1 + (1 - \lambda) A_2, \lambda B_1 + (1 - \lambda) B_2) \geq \lambda f(A_1, B_1) + (1 - \lambda) f(A_2, B_2). \quad (6.30)$$

Let's state Lieb's theorem now,

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<sup>3</sup>To prove this we introduce an auxillary system  $B$ .

**Theorem 24 (Lieb's theorem).** Let  $X$  be a matrix, and  $0 \leq t \leq 1$ . Then the function

$$f(A, B) \equiv \text{tr}(X^\dagger A^t X B^{1-t}) \quad (6.31)$$

is jointly concave in positive matrices  $A$  and  $B$ .

This theorem implies another theorem, convexity of the relative entropy.

**Theorem 25 (Convexity of the relative entropy).** The relative entropy  $S(\rho \parallel \sigma)$  is jointly convex in its arguments.

**Corollary 25.1 (Concavity of the quantum conditional entropy).** Let  $AB$  be a composite quantum system with components  $A$  and  $B$ . Then the conditional entropy  $S(A|B)$  is concave in the state  $\rho^{AB}$  of  $AB$ .

This finally leads us to strong subadditivity.

**Theorem 26.** For any trio of quantum systems  $A, B, C$  the inequalities

$$S(A) + S(B) \leq S(A, C) + S(B, C) \quad (6.32)$$

$$S(A, B, C) + S(B) \leq S(A, B) + S(B, C) \quad (6.33)$$

hold.

#### 6.4.2 Elementary applications of strong subadditivity

Strong subadditivity and related results have many useful applications in quantum information theory. Take the inequality  $S(A) + S(B) \leq S(A, C) + S(B, C)$  which is true, it also holds for Shannon entropy too, but for different reasons. For practical applications strong subadditivity is most easily applied as a rephrasing in terms of conditional and mutual entropies. Here are three major reformulations

**Theorem 3 (Reformulations of strong subadditivity)**

1. **Conditioning reduces entropy:** Suppose  $ABC$  is a composite system. Then  $S(A|B, C) \leq S(A|B)$ .
2. **Discarding quantum systems never increases mutual information:** Suppose  $ABC$  is a composite quantum system. Then  $S(A : B) \leq S(A : B, C)$ .
3. **Quantum operations never increase mutual information:** Suppose  $AB$  be a composite system and  $\mathcal{E}$  is a trace preserving quantum operation on  $B$ . Let  $S(A : B)$  be before operation,  $S(A' : B')$  be after operation, then  $S(A' : B') \leq S(A : B)$ .

We saw before that Shannon mutual information is not subadditive, thus quantum mutual information isn't subadditive either. But it turns out that conditional entropy is subadditive, i.e

$$S(A_1, A_2 | B_1, B_2) \leq S(A_1 | B_1) + S(A_2 | B_2) \quad (6.34)$$

What's more, it's subadditive in the first and second entries.

**Theorem 27 (Subadditivity of conditional entropy).** Let  $ABCD$  be a composite of four systems. Then the conditional entropy is jointly subadditive in the first two entries:

$$S(A, B|C, D) \leq S(A|C) + S(B|D) \quad (6.35)$$

Let  $ABC$  be a composite of three quantum systems. Then conditional entropy is subadditive in each of first two entries:

$$S(A, B|C) \leq S(A|C) + S(B|C) \quad (6.36)$$

$$S(A|B, C) \leq S(A|B) + S(A|C) \quad (6.37)$$

We got introduced to relative entropy like a measure between density operators. It seems intuitive that this relative entropy decreases when we ignore a part of system is ignored, i.e it makes it harder to distinguish two states of that system and thus decrease any reasonable measure of distance between them. This following theorem states it:

**Theorem 28.** Let  $\rho^{AB}$  and  $\sigma^{AB}$  be any two density matrices of a composite system  $AB$ . Then

$$S(\rho^A \| \rho^A) \leq S(\rho^{AB} \| \rho^{AB}). \quad (6.38)$$

## **Part VI**

### **Week 6**

# Chapter 7

## Quantum circuits

### 7.1 Quantum algorithms

Many interesting problems are impossible to solve on a classical computer, not because they're in principle insoluble, but because of the astronomical amount of resources they'd need. Quantum computing promises to enable feasible algorithms which would require unreasonably high resources for a solution on a classical computer.

As we've discussed before, two broad classes of quantum algorithms exist now, first is based on Shor's *Quantum fourier transform* which provides *exponential* speedup. Second is based on Grover's algorithm for performing *quantum searching* which provides *quadratic* speedup. Few main applications are listed in this figure

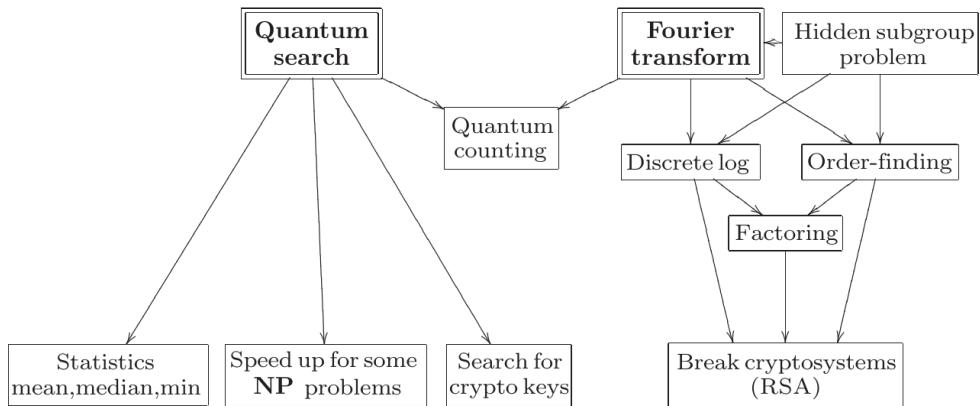


Figure 7.1: Main quantum algorithms with their applications and relationships.

### 7.2 Single qubit operations

A single qubit is the simplest quantum system. We already know few important things like norm of the state is 1, operations on it are defined by unitary matrices etc. Few important gates/ operations are Pauli matrices ( $I$ ,  $X$ ,  $Y$ ,  $Z$ ) which we know. Three other important quantum gates are Hadamard gate ( $H$ ), phase gate ( $S$ ), and  $\pi/8$  gate ( $T$ ):

$$H = \frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 1 \\ 1 & -1 \end{bmatrix}; \quad S = \begin{bmatrix} 1 & 0 \\ 0 & i \end{bmatrix}; \quad T = \begin{bmatrix} 1 & 0 \\ 0 & \exp(i\pi/4) \end{bmatrix}. \quad (7.1)$$

Few useful facts are  $H = (X + Z)/\sqrt{2}$  and  $S = T^2$ ,  $T$  is known as  $\pi/8$  gate instead of  $\pi/4$  because it's equal to a gate with diagonal  $\exp(\pm\pi/8)$  upto a global phase factor:

$$T = \exp(i\pi/8) \begin{bmatrix} \exp(-i\pi/8) & 0 \\ 0 & \exp(i\pi/8) \end{bmatrix} \quad (7.2)$$

Also, in *Bloch sphere* representation, the state  $a|0\rangle + b|1\rangle$  can be represented as a point  $(\theta, \varphi)$  where  $a = \cos(\theta/2)$  and  $b = e^{i\varphi} \sin(\theta/2)$ , and the Bloch vector is  $(\cos \varphi \sin \theta, \sin \varphi \sin \theta, \cos \theta)$ .

Pauli matrices give rise to useful matrices when exponentiated, the *rotation operators* about  $\hat{x}$ ,  $\hat{y}$  and  $\hat{z}$  axes, defined by equations

$$R_x(\theta) \equiv e^{-i\theta X/2} = \cos \frac{\theta}{2} I - i \sin \frac{\theta}{2} X = \begin{bmatrix} \cos \frac{\theta}{2} & -i \sin \frac{\theta}{2} \\ -i \sin \frac{\theta}{2} & \cos \frac{\theta}{2} \end{bmatrix} \quad (7.3)$$

$$R_y(\theta) \equiv e^{-i\theta Y/2} = \cos \frac{\theta}{2} I - i \sin \frac{\theta}{2} Y = \begin{bmatrix} \cos \frac{\theta}{2} & -\sin \frac{\theta}{2} \\ \sin \frac{\theta}{2} & \cos \frac{\theta}{2} \end{bmatrix} \quad (7.4)$$

$$R_z(\theta) \equiv e^{-i\theta Z/2} = \cos \frac{\theta}{2} I - i \sin \frac{\theta}{2} Z = \begin{bmatrix} e^{-i\theta/2} & 0 \\ 0 & e^{i\theta/2} \end{bmatrix}. \quad (7.5)$$

this is because if  $x \in \mathbb{R}$  and  $A^2 = I$ , then  $e^{iAx} = \cos xI + i \sin xA$ . If  $\hat{n} = (n_x, n_y, n_z)$  is real unit vector in three dimensions then generalize the rotation by  $\theta$  about  $\hat{n}$  by the equation

$$R_{\hat{n}}(\theta) = e^{-i\theta\hat{n}\cdot\vec{\sigma}/2} = \cos\left(\frac{\theta}{2}\right)I - \sin\left(\frac{\theta}{2}\right)(n_x X + n_y Y + n_z Z), \quad (7.6)$$

where  $\vec{\sigma}$  denotes the three component vector  $(X, Y, Z)$  of Pauli matrices. A nice fact is that any unitary operator on a single qubit can be written as  $U = e^{i\alpha} R_{\hat{n}}(\theta)$  for some  $\alpha$ ,  $\hat{n}$  and  $\theta$ . More generally

**Theorem 29 (Z-Y decomposition for a single qubit).** Suppose  $U$  is a unitary operation on a single qubit. Then there exist real numbers  $\alpha, \beta, \gamma, \delta$  such that

$$U = e^{i\alpha} R_z(\beta) R_y(\gamma) R_z(\delta) \quad (7.7)$$

In general  $z$  and  $y$  can be replaced by any non-parallel real unit vectors  $\hat{m}$  and  $\hat{n}$ . Also, this leads to a mysterious corollary, which'll be useful later.

**Corollary 29.1.** Suppose  $U$  is a unitary gate on a single qubit. Then there exist unitary operators  $A$ ,  $B$ , and  $C$  on a single qubit such that  $ABC = I$  and  $U = e^{i\alpha} AXBXC$ , where  $\alpha$  is some overall phase factor.

Here are few useful circuit identities:

$$HXH = Z; HYH = -Y; HZH = X. \quad (7.8)$$

also,  $HTH = R_x(\pi/4)$ . To recap things here are few quantum circuit symbols

Hadamard	$\boxed{H}$	$\frac{1}{\sqrt{2}} \begin{bmatrix} 1 & 1 \\ 1 & -1 \end{bmatrix}$
Pauli-X	$\boxed{X}$	$\begin{bmatrix} 0 & 1 \\ 1 & 0 \end{bmatrix}$
Pauli-Y	$\boxed{Y}$	$\begin{bmatrix} 0 & -i \\ i & 0 \end{bmatrix}$
Pauli-Z	$\boxed{Z}$	$\begin{bmatrix} 1 & 0 \\ 0 & -1 \end{bmatrix}$
Phase	$\boxed{S}$	$\begin{bmatrix} 1 & 0 \\ 0 & i \end{bmatrix}$
$\pi/8$	$\boxed{T}$	$\begin{bmatrix} 1 & 0 \\ 0 & e^{i\pi/4} \end{bmatrix}$

Figure 7.2: Names, symbols and unitary matrices for commonly used quantum gates.

### 7.3 Controlled operations

If  $A$  is true, then do  $B'$  is the most useful and basic controlled operation in both classical and quantum computation. An example is CNOT gate in quantum computation, which does  $|c\rangle|t\rangle \rightarrow |c\rangle|t \oplus c\rangle$  where  $|c\rangle$  is control qubit and  $|t\rangle$  the target qubit. If  $|c\rangle$  is in state  $|0\rangle$  nothing happens to target qubits, else if it's in  $|1\rangle$   $|t\rangle$  is flipped. A general thing for this is *controlled-U* gate, which does  $|c\rangle|t\rangle \rightarrow |c\rangle U^c|t\rangle$ , where  $U$  is a unitary operation. In the computational basis  $|\text{control}, \text{target}\rangle$ , the matrix representation of CNOT gate is

$$\begin{bmatrix} 1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & 1 & 0 \end{bmatrix}. \quad (7.9)$$

To develop *controlled-U* gate for arbitrary  $U$ , we'll use  $U = e^{i\alpha}AXBXC$ . To apply the controlled phase shift  $\exp(i\alpha)$ , we use a circuit containing a single qubit gate as shown

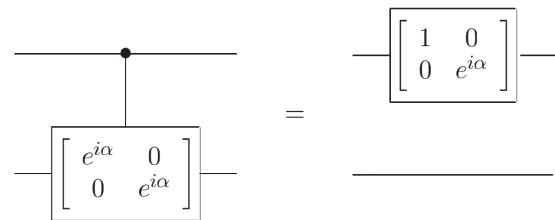


Figure 7.3: Controlled phase shift gate and an equivalent circuit for two qubits.

To complete the gate we use the circuit 7.4. This works because if control qubit is set to  $|0\rangle$  then nothing happens at CNOT gates and  $ABC = I$  is applied which does nothing, if control qubit is  $|1\rangle$  then  $e^{i\alpha}AXBXC$  is applied, thus  $U$  is applied.

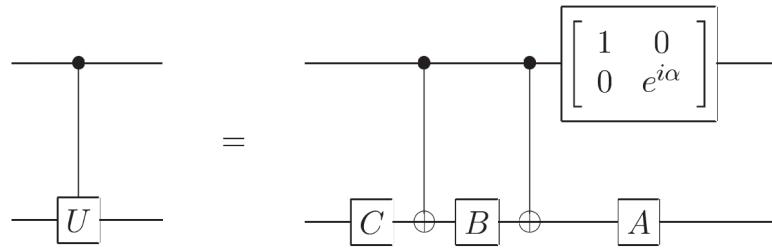


Figure 7.4: *controlled-U* circuit,  $\alpha, A, B, C$  satisfy  $U = e^{i\alpha} AXBXC$  and  $ABC = I$

To use conditioning on multiple-qubits, suppose we have  $n + k$  qubits, and  $U$  is a  $k$  bit operator. Then the controlled operation  $C^n(U)$  by the equation

$$C^n(U) |x_1 x_2 \dots x_n\rangle |\psi\rangle = |x_1 x_2 \dots x_n\rangle U^{x_1 x_2 \dots x_n} |\psi\rangle \quad (7.10)$$

where  $x_1 x_2 \dots x_n$  means product of those bits. It means operation  $U$  is applied on last  $k$  qubits if all of the first  $n$  qubits are  $|1\rangle$ . We'd assume  $k = 1$ , for  $k \geq 2$  we don't yet know how to perform  $k$  arbitrary operations at once.

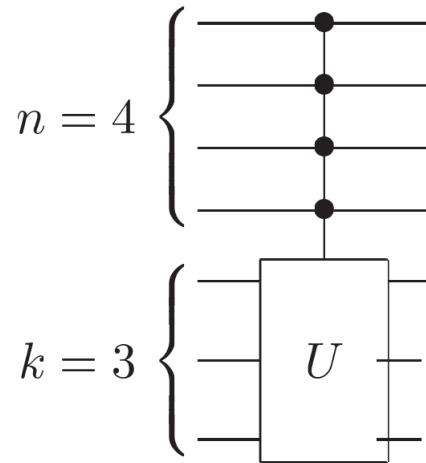


Figure 7.5:  $C^n(U)$  operation for  $n = 4$  and  $k = 3$ .

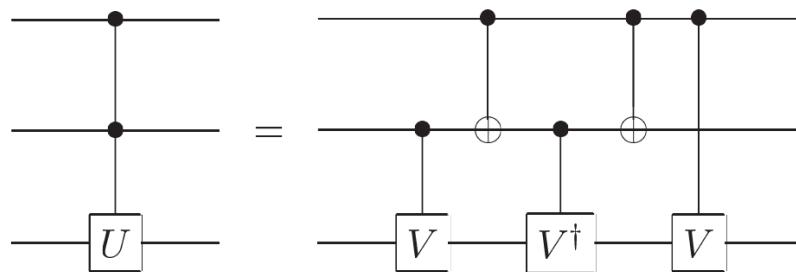


Figure 7.6:  $C^2(U)$  gate.  $V$  is a unitary operator satisfying  $V^2 = U$ . When  $V \equiv (1 - i)(1 + iX)/2$  represents the Toffoli gate.

Remember that when  $U = X$ ,  $C^2(U)$  is known as the *Toffoli gate*.

How  $C^n(U)$  is implemented is interesting, first we take the first  $n$  qubits as  $|c_1\rangle, |c_2\rangle \dots |c_n\rangle$ , we use  $(n-1)$  intermediate working qubits which are all initially and finally  $|0\rangle$ , which store  $|c_1c_2\rangle, |c_1c_2c_3\rangle \dots |c_1c_2 \dots c_n\rangle$ . This is done using toffoli gates which is applied between  $|c_1\rangle$  and  $|c_2\rangle$  with the output to first working qubit, the next toffoli gate is applied between first working qubit and  $|c_3\rangle$  which gives  $|c_1c_2c_3\rangle$  and so on, thus giving  $|c_1c_2 \dots c_n\rangle$ .

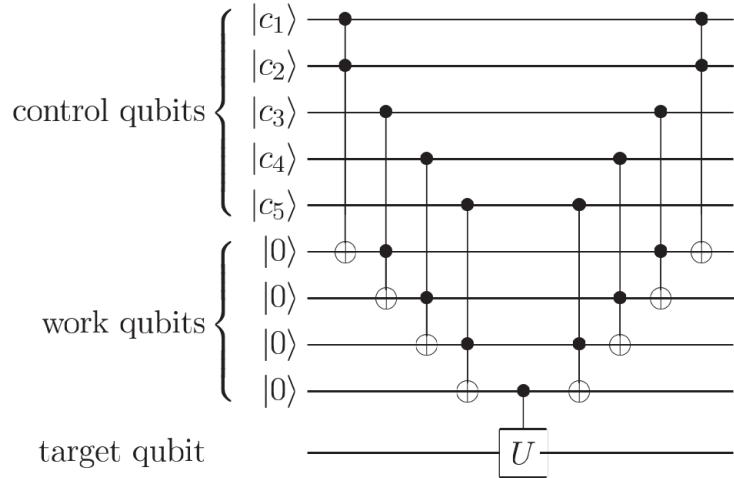


Figure 7.7: Implementation of  $C^n(U)$  circuit.

We are considering conditionals being 1 till now, but sometimes it's useful to keep the conditional as 0 for which we introduce a notation, we keep an open circle (without darkening it) whereas we used to have a darkened circle till now.

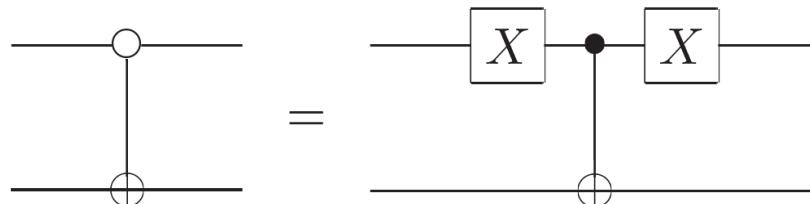


Figure 7.8: CNOT GATE, WITH THE CONDITIONAL BEING THAT FIRST QUBIT IS SET TO 0.

## 7.4 Measurement

As we've shown before measurements will be shown by a 'meter' symbol, especially projective measurements. There are two major principles worth emphasizing early, first is that classically conditioned operations can be replaced by quantum controlled operations:

**Principle of deferred measurement:** Measurements can always be moved from an intermediate state of a quantum circuit to the end of the circuit; if the measurement results are used at any stage, then the classically conditioned operations can be replaced by conditional quantum operations.

**Principle of implicit measurement:** Without loss of generality, any unterminated quantum wires (qubits which are not measured) at the end of a quantum circuit may be assumed to be measured.

## 7.5 Universal quantum gates

If a set of gates can compute any arbitrary classical function, they are said to be *universal* for classical computation. Toffoli gate is universal for classical computation, hence quantum computation subsumes classical computation.

### 7.5.1 Two level unitary gates are universal

We'll see how to decompose a  $d$ -dimensional unitary matrix  $U$  into a product of *two-level unitary matrices*, which act non-trivially on only two or fewer components. We'll understand what're these by going forward. Let's take an example  $U$  where it's  $3 \times 3$ ,

$$U = \begin{bmatrix} a & d & g \\ b & e & h \\ c & f & i \end{bmatrix}. \quad (7.11)$$

We can find two-level unitary matrices such that

$$U_3 U_2 U_1 U = I \quad (7.12)$$

or,  $U = U_1^\dagger U_2^\dagger U_3^\dagger$ . More generally, for  $d$ -dimensional space,  $U$  can be written as

$$U = V_1 \cdots V_k \quad (7.13)$$

where the matrices  $V_i$  are unitary matrices and  $k \leq (d-1) + (d-2) + \cdots + 1 = d(d-1)/2$ . This follows that any unitary operator on  $n$  qubits can be written as a product of  $2^{n-1}(2^n - 1)$  two-level unitary matrices.

### 7.5.2 A discrete set of universal operators

#### Approximating unitary operators

A *discrete* set of gates can't describe any unitary operator, since the latter are continuous. Suppose we want to implement operator  $U$  but what we actually implement is  $V$ , the *error* when  $V$  is implemented instead of  $U$  is

$$E(U, V) \equiv \max_{|\psi\rangle} \|(U - V)|\psi\rangle\| \quad (7.14)$$

maximum is over all normalized states  $|\psi\rangle$  of state space.

#### Universality of Hadamard + phase + CNOT + $\pi/8$ gates

We're going to study two different discrete set of gates, both of which are universal. *Standard set* consists of Hadamard, CNOT and  $\pi/8$  gates. The second set of gates contains Hadamard, phase, CNOT and Toffoli gates. All of these gates can be done fault tolerantly.

It can be shown that using these gates, any rotation can be done up to desired accuracy. Given any single qubit operator  $U$  and any  $\epsilon > 0$  it's possible to approximate  $U$  within  $\epsilon$  using a circuit composed of Hadamard and  $\pi/8$  gates alone. This simulation is efficient, because the overhead required to perform the simulation is polynomial in  $\log(m/\epsilon)$ , where  $m$  is number of gates in circuit and  $\epsilon$  is desired accuracy.

### 7.5.3 Approximating arbitrary unitary gates is generically hard

There are states of  $n$  qubits which take  $\Omega(2^n \log(1/\epsilon) / \log(n))$  operations to approximate to within a distance  $\epsilon$ . This is exponential in  $n$ , and thus is 'difficult'. This implies that there are unitary transformations  $U$  on  $n$  qubits which take  $\Omega(2^n \log(1/\epsilon) / \log(n))$  operations to approximate by a quantum circuit implementing an operation  $V$  such that  $E(U, V) \leq \epsilon$ .

## 7.6 Simulation of quantum systems

To simulate, solutions of differential equations capturing physical laws governing the dynamical behaviour of the system. For example the Newton's law,

$$\frac{d}{dt} \left( m \frac{dx}{dt} \right) = F, \quad (7.15)$$

Poisson's equation,

$$-\vec{\nabla} \cdot (k \vec{\nabla} \vec{u}) = \vec{Q}, \quad (7.16)$$

the electromagnetic vector wave equation,

$$\vec{\nabla} \cdot \vec{\nabla} \vec{E} = \epsilon_0 \mu_0 \frac{\partial^2 \vec{E}}{\partial t^2}, \quad (7.17)$$

and the diffusion equation,

$$\vec{\nabla}^2 \psi = \frac{1}{a^2} \frac{\partial \psi}{\partial t}, \quad (7.18)$$

Solutions for these are obtained by using *approximation*. However simulation of quantum systems by classical computers is inefficient. Many simple quantum systems are governed by

$$i\hbar \frac{d}{dt} |\psi\rangle = H |\psi\rangle. \quad (7.19)$$

For a Hamiltonian dealing with particles in space, this reduces to

$$i \frac{\partial}{\partial x} \psi(x) = \left[ -\frac{1}{2m} \frac{\partial^2}{\partial x^2} + V(x) \right] \psi(x), \quad (7.20)$$

using a convention known as the position representation  $\langle x|\psi\rangle = \psi(x)$ . The key challenge is there *exponential* number of differential equations to be solved. For  $n$  qubits,  $2^n$  equations.

Quantum computers can efficiently simulate quantum systems for which there is no known efficient classical simulation. This is similar to the fact that any quantum circuit can be designed from a discrete set of gates. Ofcourse, there are few operators which can't be efficiently approximated, similarly there are few systems with Hamiltonians which can't be efficiently simulated on a quantum computer.

### 7.6.1 The quantum simulation algorithm

Similar to how we use  $y(t + \Delta t) \approx y(t) + f(y)\Delta t$  to solve  $dy/dt = f(y)$ . To solve  $i\hbar |\psi\rangle / dt = H |\psi\rangle$ , which for time independent  $H$  is,

$$|\psi(t)\rangle = e^{-iHt} |\psi(0)\rangle. \quad (7.21)$$

Similarly, exponentiating  $H$  would be very inefficient, thus we use  $|\psi(t + \Delta t)\rangle \approx (I - iH\Delta t)|\psi(t)\rangle$ . However this isn't very satisfactory. But this can be improved by writing it as a sum of Hamiltonians acting on fewer qubits.

The heart of quantum simulation algorithms is the following asymptotic approximation theorem:

**Theorem 30 (Trotter formula).** *Let  $A$  and  $B$  be Hermitian operators. Then for any real  $t$ ,*

$$\lim_{n \rightarrow \infty} (e^{iAt/n} e^{iBt/n})^n = e^{i(A+B)t}. \quad (7.22)$$

Here's an overview of quantum simulation algorithm,

**Algorithm 2:** Quantum simulation

**Inputs :** (1) A Hamiltonian  $H = \sum_k H_k$  acting on an  $N$ -dimensional system, where each  $H_k$  acts on a small subsystem of size independent of  $N$ , (2) an initial state  $|\psi_0\rangle$ , of the system at  $t = 0$ , (3) a positive, non-zero accuracy  $\delta$ , and (4) a time  $t_f$  at which the evolved state is desired.

**Outputs :** A state  $|\tilde{\psi}(t_f)\rangle$  such that  $|\langle \tilde{\psi}(t_f) | e^{-iHt_f} |\psi_0\rangle|^2 \geq 1 - \delta$ .

**Runtime:**  $O(\text{poly}(1/\delta))$  operations.

**Procedure**

Choose a representation such that the state  $|\tilde{\psi}\rangle$  of  $n = \text{poly}(\log N)$  qubits approximates the system and the operators  $e^{-iH_k\Delta t}$  have efficient quantum circuit approximations. Select an approximation method and  $\Delta t$  such that the expected error is acceptable (and  $j\Delta t = t_f$  for an integer  $j$ ), construct the corresponding quantum circuit  $U_{\Delta t}$  for the iterative step, and do:

$|\tilde{\psi}_0\rangle \leftarrow |\psi_0\rangle; j = 0$  // initialize state

$\rightarrow |\tilde{\psi}_{j+1}\rangle = U_{\Delta t} |\tilde{\psi}_j\rangle$  // iterative update

$\rightarrow j = j + 1$ ; goto step 2 until  $j\Delta t \geq t_f$  // loop

$\rightarrow |\tilde{\psi}(t_f)\rangle = |\tilde{\psi}_j\rangle$  // final result

## **Part VII**

### **Week 7**

# Chapter 8

## The quantum Fourier transform and its applications

We'll study *quantum Fourier transform*, which is a class of problems quantum computers are more efficient than classical computers.

### 8.1 The quantum Fourier transform

The *discrete Fourier transform* takes a vector of complex numbers  $x_0, \dots, x_{N-1}$  of complex numbers and outputs a same length vector of complex numbers  $y_0, \dots, y_{N-1}$  given by

$$y_k \equiv \frac{1}{\sqrt{N}} \sum_{j=0}^{N-1} x_j e^{2\pi i j k / N}. \quad (8.1)$$

The *quantum Fourier transform* is exactly the same thing but with a different convention. The quantum Fourier transform on an orthonormal basis  $|0\rangle, \dots, |N-1\rangle$  is a linear operator which does

$$|j\rangle \longrightarrow \frac{1}{\sqrt{N}} \sum_{j=0}^{N-1} |k\rangle e^{2\pi i j k / N}. \quad (8.2)$$

which corresponds to

$$\sum_{j=0}^{N-1} x_j |j\rangle \longrightarrow \sum_{k=0}^{N-1} y_k |k\rangle \quad (8.3)$$

where  $y_k$  are discrete Fourier transforms of  $x_j$ . It's not obvious but this is a unitary transform, and thus can be implemented by a quantum computer efficiently.

If  $N = 2^n$  and we can write  $|j\rangle$  in its binary representation as  $|j_1, j_2, \dots, j_n\rangle$  and we also write  $0.j_l j_{l+1} \dots j_m$  as the *binary fraction*  $j_l/2 + j_{l+1}/4 + \dots + j_m/2^{m-l+1}$  then the quantum Fourier transform in its *product notation* is

$$|j_1, \dots, j_n\rangle \longrightarrow \frac{(|0\rangle + e^{2\pi i 0 \cdot j_n} |1\rangle) (|0\rangle + e^{2\pi i 0 \cdot j_{n-1} j_n} |1\rangle) \dots (|0\rangle + e^{2\pi i 0 \cdot j_1 j_2 \dots j_n} |1\rangle)}{2^{n/2}}. \quad (8.4)$$

The evidence of the above product representation is as followed:

$$|j\rangle \rightarrow \frac{1}{2^{n/2}} \sum_{k=0}^{2^n-1} e^{2\pi i j k / 2^n} |k\rangle \quad (8.5)$$

$$= \frac{1}{2^{n/2}} \sum_{k_1=0}^1 \dots \sum_{k_n=0}^1 e^{2\pi i j (\sum_{l=1}^n k_l 2^{-l}) / 2^n} |k_1 \dots k_n\rangle \quad (8.6)$$

$$= \frac{1}{2^{n/2}} \sum_{k_1=0}^1 \dots \sum_{k_n=0}^1 \bigotimes_{l=1}^n e^{2\pi i j k_l 2^{-l}} |k_l\rangle \quad (8.7)$$

$$= \frac{1}{2^{n/2}} \bigotimes_{l=1}^n \left[ \sum_{k_l=0}^1 e^{2\pi i j k_l 2^{-l}} |k_l\rangle \right] \quad (8.8)$$

$$= \frac{1}{2^{n/2}} \bigotimes_{l=1}^n \left[ |0\rangle + e^{2\pi i j 2^{-l}} |1\rangle \right] \quad (8.9)$$

$$= \frac{(|0\rangle + e^{2\pi i 0 \cdot j_n} |1\rangle) (|0\rangle + e^{2\pi i 0 \cdot j_{n-1} j_n} |1\rangle) \dots (|0\rangle + e^{2\pi i 0 \cdot j_1 j_2 \dots j_n} |1\rangle)}{2^{n/2}}. \quad (8.10)$$

Here's a circuit for implementing quantum Fourier transform

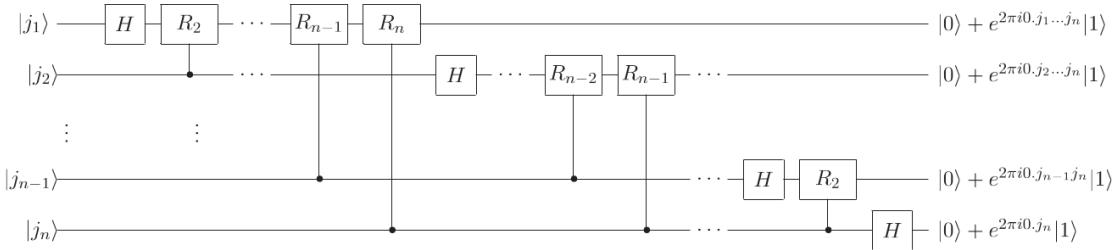


Figure 8.1: Here  $R_k = \begin{bmatrix} 1 & 0 \\ 0 & e^{2\pi i / 2^k} \end{bmatrix}$ . This circuit efficiently computes quantum fourier transform.

In this, a total of  $n(n-1)/2$  gates are required, and atmost  $n/2$  swaps are required at the end. Thus providing a  $\Theta(n^2)$  algorithm. Unfortunately, amplitudes in a quantum computer cannot be directly accessed by measurement. Thus, there's no way of determining Fourier transformed amplitudes of the original state.

## 8.2 Phase estimation