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Holonomic Optimal Control for Qudits

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Abstract

Non-adiabatic Holonomic Quantum Computation (NHQC) is a method used to implement quantum gates with holonomies that arise solely from non-Abelian geometric phases. Due to high noise tolerance these phases can be used to construct resilient quantum gates. By using dark paths we show how to implement quantum gates for higher dimensional computation elements, qudits, instead of the conventional qubits. This gives higher parameter control compared to earlier implementations. We present a scheme that generalizes and achieves single-qudit universality using controllable high fidelity gates by including an auxiliary state. An explicit example is shown for the Qutrit. The scaling is linear in dimension and we show how any diagonal qudit gate can be implemented efficiently in any dimension.

Sammanfattning

Kvantmekaniska periodiska system uppvisar, förutom den väletablerade dynamiska fasen i tidsutvecklingen, även en geometrisk fas. Eftersom denna fas är oberoende av systemets dynamik kan den användas för att skapa fel-toleranta kvantmekaniska grindar. Detta är grunden för s.k. Icke-adiabatic Holonomiska Kvantgrindar. Genom att utnyttja vägar i parameterrummet som inte bidrar med någon dynamisk utveckling så implementerar vi hög dimensionella beräkningselement, 'qudits', istället för de mer konventionella kvantbitarna. Detta resulterar i högre kontroll över parameterar än vissa tidigare metoder. Vi visar hur detta schema kan generaliseras till högre dimensioner och hur man kan uppnå universalitet med dessa hög-dimensionella och fel-toleranta kvantgrindar. Ett exempel ges för tre dimensioner och används som grund som sedan generaliseras. Schemat har linjär skalning i dimension och diagonala qudit-grindar kan effektivt bli implementerade i höga dimensioner.

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1 Introduction

The emerging field of quantum technology has many promising applications, one of them is quantum computation (QC), which currently is a very active area of research. Quantum computers make use of quantum mechanical effects such as superposition, entanglement, and interference to design powerful algorithms. These algorithms could be used to solve some hard problems which would not be classically possible, such as efficient prime-number factoring [1]. It can also be used to reduce the time complexity of some commonly used algorithms [2]. Current quantum computers are very susceptible to decoherence and noise, and thus will not have any commercial use any time soon, but stand as an important proof of concept. The most common model for quantum computation is the circuit model, which is analogous to the classical circuits used for classical computers. Gates are replaced by unitary transformations (quantum gates) and bits by qubits. To achieve the computational advantage it is important to construct robust, noise-resilient quantum gates. A good candidate for this is holonomic quantum computation [3, 4], which is based on the Berry phase [5] and its non-abelian and/or non-adiabatic generalizations [6, 7, 8]. These methods are only dependent on the geometry of the system and thus are resilient to local errors in the dynamical evolution.

The idea that elements of computation should be limited to two-dimensional qubits is sort of an arbitrary choice that most likely rose out of convenience due to binary logic. So why binary logic? It is simply the easiest non-trivial example, in binary things can be either 1 or 0, True or False, **on** or **off**. Due to its simplicity, it is no wonder that this is how the first computer was designed. But are we limited to qubits? As early as 1840 a mechanical ternary (three-valued logic) calculation device was built by Thomas Fowler [9], and in 1958 the first electronic ternary computer was built by the Soviet Union [10]. Even though it had many advantages over the binary computer it never saw the same widespread success. There is nothing in theory that forbids a higher dimensional computational basis, even more so when it comes to quantum computers, where the implementation of the elements of computation already surpasses the simplicity of **on** and **off**. There are already promising qudit results that show potential [11, 12, 13], and in the review article [14] a good overview of the field is given and further research into the topic is encouraged.

In this report we will show how to find a new scheme to implement qudits which could be more efficient than some current ones by making use of dark paths for increased parameter control and auxiliary states for increased fidelity. We do this by generalizing the idea of the scheme from [15]. The report is structured as follows. The background/theory section is split into two parts where the first part serves as a quick introduction to the most important aspects of quantum mechanics as well as the commonly used notation. Then follows a part more concerned with quantum computation, quantum information, and some of the more advanced quantum mechanical concepts that those are built upon. Then the main results are shown, first an explicit example of the qutrit and then how it generalizes in higher dimension. The report ends with conclusions and a brief outlook.

2 Theoretical background

2.1 Basic quantum mechanics

This part offers a quick introduction to the necessary quantum mechanics for readers who are not familiar with the subject. Quantum mechanics is the framework developed to handle modern physics and is famously somewhat hard to understand intuitively; a good baseline understanding is however obtainable with knowledge of linear algebra combined with the postulates of QM. But to understand the contents of this project it should be enough with linear algebra and some physical interpretation. The section focuses mostly on familiarising the reader with the notation used in quantum mechanics and some fundamental ideas. I suggest chapters 1 and 2 of [16] for a more complete introduction.

2.1.1 Quantum states and Dirac notation

First, let us define what is meant by the term **state**. In classical physics, a state is defined by the position and momentum of all its constituents. An example is a system of N particles, the state of which would be given by $\{(\vec{x}, \vec{p})_i\}_{i=1}^N$, where $\vec{x}, \vec{p} \in \mathbb{R}^3$ are the position and momentum in 3 dimensions. In Quantum Mechanics (QM) it is more subtle than this since exact information of the system cannot be obtained in the same way. A state is represented by a normalized vector in a **Hilbert space**, \mathcal{H} , a real or complex inner product space. For $u, v \in \mathcal{H}$ the inner product has the following properties:

1. The inner product is conjugate symmetric,

$$\langle u, v \rangle = \overline{\langle v, u \rangle} \in \mathbb{C}$$

2. The inner product is linear in the first argument, for constants $a, b \in \mathbb{C}$,

$$\langle au_1 + bu_2, v \rangle = a\langle u_1, v \rangle + b\langle u_2, v \rangle$$

3. The inner product is positive definite,

$$\langle u, u \rangle = 0 \iff u = 0$$

These properties can be combined to find some other useful facts. Combining the 1st and 2nd property, one obtains

$$\langle v, au_1 + bu_2 \rangle = \overline{\langle au_1 + bu_2, v \rangle} = a^* \overline{\langle u_1, v \rangle} + b^* \overline{\langle u_2, v \rangle} = a^* \langle v, u_1 \rangle + b^* \langle v, u_2 \rangle \quad (1)$$

the inner product is anti-linear in the second term. The 1st and 3rd property yield

$$\langle u, u \rangle = \overline{\langle u, u \rangle} \implies \text{Im}(\langle u, u \rangle) = 0, \quad (2)$$

or in words, the inner product of two identical vectors is a real number.

A quantum state is then represented by a normalized vector v , $\langle v, v \rangle = 1$, in a Hilbert space \mathcal{H} . Now a property of vector spaces is that any vector can be multiplied by a matrix, and the resulting vector will be a new vector in the same vector space. This is what is meant when a quantum state is **acted** upon. For a matrix A , we have that

$$\mathcal{H} \ni v \xrightarrow{A} Av = v' \in \mathcal{H}. \quad (3)$$

That is, acting on a state alters it in various ways. Now let us go from this linear algebra notation to the Dirac notation commonly used in quantum mechanics, also known as bra-ket notation. Vectors are replaced by **kets**, $|\rangle$, or **bras**, $\langle|$,

$$v \mapsto |\psi\rangle$$

$$v^\dagger \mapsto \langle\psi|$$

and matrices are replaced by linear operators

$$A \mapsto \hat{A}.$$

The same rules apply to these as for the usual vectors. The labels inside the brackets do not in themselves have any meanings and are in some sense only that, labels, but more often than not it is used to represent some property of the state. With this notation the inner product between to states $|\psi\rangle, |\varphi\rangle$ is written and defined as

$$\langle\psi|\varphi\rangle = a_1^*b_1 + a_2^*b_2 + \cdots + a_n^*b_n = \begin{pmatrix} a_1^* & a_2^* & \cdots & a_n^* \end{pmatrix} \begin{pmatrix} b_1 \\ b_2 \\ \vdots \\ b_n \end{pmatrix}. \quad (4)$$

From the properties defined earlier the relations $(|\psi\rangle)^\dagger = \langle\psi|$ and $(\langle\psi|\varphi\rangle)^\dagger = \langle\varphi|\psi\rangle$, suggests another way to define the states simply as

$$|\psi\rangle = \begin{pmatrix} a_1 \\ a_2 \\ \vdots \\ a_n \end{pmatrix}, \quad \langle\psi| = (|\psi\rangle)^\dagger = (a_1^*, a_2^*, \dots, a_n^*), \quad a_1, a_2, \dots, a_n \in \mathbb{C}, \quad (5)$$

which is nothing more than a complex vector. A common way to represent kets is as a linear combination of eigenkets, which corresponds to the set of eigenvectors. Thus, any ket can be written as a linear combination of these kets if they span the Hilbert space

$$|\Psi\rangle = \sum_{\psi} c_{\psi} |\psi\rangle. \quad (6)$$

Following are some example calculations to see how to use Dirac-notation. Let us consider what the effect of the Pauli-Z operator of a spin- $\frac{1}{2}$ particle is

a) on the up position $|\psi\rangle = |\downarrow\rangle$

- b)** on an arbitrary superposition of up and down $|\psi\rangle = \alpha|\uparrow\rangle + \beta|\downarrow\rangle$, $|\alpha|^2 + |\beta|^2 = 1$, $\alpha, \beta \in \mathbb{C}$

Let us first show how this would be done using the matrix representation and then using Dirac notation. Here the kets $\{|\uparrow\rangle \doteq \begin{pmatrix} 1 \\ 0 \end{pmatrix}, |\downarrow\rangle \doteq \begin{pmatrix} 0 \\ 1 \end{pmatrix}\}$ span the Hilbert space of the particle, they also form an orthonormal basis. The Pauli-Z operator in the matrix representation in this basis is $\sigma_z = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}$. Then we have that

- a)** We simply write out the matrices and vectors and get that

$$\sigma_z |\psi\rangle \doteq \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} \begin{pmatrix} 0 \\ 1 \end{pmatrix} = \begin{pmatrix} 0 \\ -1 \end{pmatrix} \doteq -|\downarrow\rangle \quad (7)$$

- b)** Same as before but with a superposition (linear combination) up and down

$$\begin{aligned} \sigma_z |\psi\rangle &\doteq \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} \left[\alpha \begin{pmatrix} 1 \\ 0 \end{pmatrix} + \beta \begin{pmatrix} 0 \\ 1 \end{pmatrix} \right] \\ &= \begin{pmatrix} \alpha \\ -\beta \end{pmatrix} \\ &= \alpha \begin{pmatrix} 1 \\ 0 \end{pmatrix} - \beta \begin{pmatrix} 0 \\ 1 \end{pmatrix} \\ &= \alpha |\uparrow\rangle - \beta |\downarrow\rangle \end{aligned} \quad (8)$$

The matrix representation works great for small Hilbert spaces since manipulating small matrices by hand is not very time consuming. Consider now the same calculation using Dirac notation, here the inner product will be a key aspect, and making use of the orthonormality of $|\uparrow\rangle, |\downarrow\rangle$ it follows that $\langle\uparrow|\uparrow\rangle = \langle\downarrow|\downarrow\rangle = 1$ and $\langle\uparrow|\downarrow\rangle = \langle\downarrow|\uparrow\rangle = 0$. With Dirac notation operators can be expressed as outer products of bras and kets, thus Pauli-Z can be written as $\sigma_z = |\uparrow\rangle\langle\uparrow| - |\downarrow\rangle\langle\downarrow|$.

- a)**

$$\begin{aligned} \sigma_z |\psi\rangle &= (|\uparrow\rangle\langle\uparrow| - |\downarrow\rangle\langle\downarrow|) |\downarrow\rangle \\ &= |\uparrow\rangle \underbrace{\langle\uparrow|\downarrow\rangle}_{=0} - |\downarrow\rangle \underbrace{\langle\downarrow|\downarrow\rangle}_{=1} \\ &= 0 |\uparrow\rangle - 1 |\downarrow\rangle \\ &= -|\downarrow\rangle. \end{aligned} \quad (9)$$

- b)**

$$\begin{aligned} \sigma_z |\psi\rangle &= (|\uparrow\rangle\langle\uparrow| - |\downarrow\rangle\langle\downarrow|) (\alpha |\uparrow\rangle + \beta |\downarrow\rangle) \\ &= \alpha |\uparrow\rangle \underbrace{\langle\uparrow|\uparrow\rangle}_{=1} - \alpha |\downarrow\rangle \underbrace{\langle\uparrow|\downarrow\rangle}_{=0} + \beta |\uparrow\rangle \underbrace{\langle\downarrow|\uparrow\rangle}_{=0} - \beta |\downarrow\rangle \underbrace{\langle\downarrow|\downarrow\rangle}_{=1} \\ &= \alpha |\uparrow\rangle - \beta |\downarrow\rangle. \end{aligned} \quad (10)$$

Note that the final answer is the same for both cases, but using Dirac notation there is need to bother with explicit vectors and matrices, that is one of the strong advantages of the Dirac notation. These calculations are quite analogous to how a quantum gate acting on a qubit works which will be discussed further in the background. In Dirac notation any general matrix A can be written as $A = \sum_{ij} A_{ij} |i\rangle \langle j|$. Measuring a quantum state $|\psi\rangle = \sum_n c_n |n\rangle$, $c_n \in \mathbb{C}$ collapses it into a single base ket of the measured basis

$$|\psi\rangle \xrightarrow{\text{Measurement}} |n\rangle, \quad (11)$$

with probability $|c_n|^2$.

2.1.2 Time evolution and the Schrödinger equation

Say we have a time-dependent state $|\psi(t)\rangle$. How does the system evolve when time goes from t_0 to some later time $t_1 > t_0$,

$$|\psi(t_0)\rangle \xrightarrow{\text{Time evolution}} |\psi(t_1)\rangle \quad (12)$$

We can write this as an operator equation $U(t_1, t_0) |\psi(t_0)\rangle = |\psi(t_1)\rangle$, where U is called the time evolution operator. To determine the form of U let put some constraints on U . To preserve probability the operator must be unitary, $U^\dagger U = U U^\dagger = 1$. The second property that needs to be fulfilled is composition, $U(t_2, t_0) = U(t_2, t_1) U(t_1, t_0)$, which is to say that applying two consecutive time transformations, first between $t_0 \rightarrow t_1$ and then $t_1 \rightarrow t_2$ is equivalent to applying one transformation between $t_0 \rightarrow t_2$. Proceeding let us define how the time evolution operator would act under an infinitesimal change dt , let's propose an operator on the form

$$U(t_0 + dt, t_0) = \mathbf{1} - i \Omega dt \quad (13)$$

with Ω being Hermitian, $\Omega = \Omega^\dagger$. Ω can be time dependent, and then must be evaluated at t_0 . As dt approaches 0 it reduces to an identity operator, which is appropriate. Let's check the two other properties

1. Unitarity

$$\begin{aligned} U(t_0 + dt, t_0) U(t_0 + dt, t_0)^\dagger &= (\mathbf{1} - i \Omega dt) (\mathbf{1} + i \Omega dt) \\ &= \mathbf{1} + i \Omega dt - i \Omega dt + \mathcal{O}(dt^2) \\ &\simeq \mathbf{1} \end{aligned} \quad (14)$$

2. Composition

$$\begin{aligned} U(t_0 + dt_1 + dt_2, t_0 + dt_1) U(t_0 + dt_1, t_0) &= (\mathbf{1} - i \Omega dt_2) (\mathbf{1} - i \Omega dt_1) \\ &= \mathbf{1} - i \Omega dt_1 - i \Omega dt_2 + \mathcal{O}(dt^2) \\ &\simeq \mathbf{1} - i \Omega (dt_1 + dt_2) \\ &= U(t_0 + dt_1 + dt_2, t_0). \end{aligned} \quad (15)$$

From the operator Ω the Hamiltonian, H , is defined, $\Omega = \frac{H}{\hbar}$, where \hbar is Planck's reduced constant. Thus by definition the Hamiltonian will always be Hermitian by

definition. So for an infinitesimal time shift the operator has the form $U(t_0 + dt, t_0) = 1 - i \frac{H}{\hbar} dt$. Time evolution is closely tied to one of Quantum Mechanics postulates, that is that a closed quantum state evolves according to the Schrödinger equation, Equation 16, here it becomes more obvious why it happens that Ω was related to the Hamiltonian. It is common to set $\hbar = 1$, as it is just a numerical shift, so let's do that, which simplifies both the Schrödinger equation and the time evolution operator U .

$$i \frac{\partial}{\partial t} |\psi(t)\rangle = H |\psi(t)\rangle \quad (16)$$

Now it is clear that the time evolution operator is related to the Schrödinger equation in some way, notice that the solution $|\psi(t)\rangle$ can be written in terms of our time evolution operator as $|\psi(t)\rangle = U(t, t_0) |\psi(t_0)\rangle$. So given an initial condition we have solved the Schrödinger equation if we can determine $U(t, t_0)$. To find the time evolution operator for any time t , we could successively repeat the infinitesimal operator until the desired time is reached $t_0 \mapsto t_0 + dt \mapsto t_0 + 2dt \mapsto \dots \mapsto t - dt \mapsto t$. For the case of a time independent Hamiltonian, H , it could be written as

$$\lim_{N \rightarrow \infty} \left(1 - i H \frac{(t - t_0)}{N} \right)^N = \exp(-i H(t - t_0)) \quad (17)$$

For simplicity let us assume that $t_0 = 0$, then the time evolution operator for a finite time for a time independent Hamiltonian could be written as $U(t, 0) = e^{-i H t}$. If the Hamiltonian has time dependence but it commutes with itself at different times, then an integral can be introduced in the exponential, $U(t, 0) = e^{-i \int_0^t H(t') dt'}$, if the Hamiltonian does not commute at different times it becomes a bit more involved and will not be covered here. But the formal solution can be written as $U(t, 0) = \mathcal{T} e^{-i \int_0^t H(t') dt'}$, where \mathcal{T} is time-ordering. This is referred to as the dynamical phase and is the most common type of time evolution. In addition to the dynamic phase there exist a geometric phase associated with some quantum systems and was first studied in periodic systems [5]. This is discussed further in Section Geometric quantum computation.

To conclude, finding the time evolution operator corresponds to finding the solution to the Schrödinger Equation, which governs the time evolution of quantum systems.

2.2 Quantum computation and quantum information theory

2.2.1 The qubit

A classical bit is a two-state (binary) system, which means it can occupy either one of two states, 0 or 1. So with n bits, there are 2^n possible states.

The general way to define a qubit would be any two dimensional quantum system, meaning that the Hilbert space is of dimension 2, $\dim(\mathcal{H}) = 2$. Where qubits are on the form given in Equation 18, any normalized superposition (linear combination) of the basis-kets $|0\rangle, |1\rangle$.

$$|\psi\rangle = \alpha |0\rangle + \beta |1\rangle, \alpha, \beta \in \mathbb{C}, |\alpha|^2 + |\beta|^2 = 1. \quad (18)$$

Most of the time it is better to work with the abstract concept of qubits without focusing on the physical implementation, however, we see a lot of similarities with spin- $\frac{1}{2}$ particles $\{|\downarrow\rangle, |\uparrow\rangle\}$ or energy levels in a trapped ion $\{|g\rangle, |e\rangle\}$. That is because abstractly those systems can be represented by the exact same mathematical representation in quantum mechanics, they are all 2-dimensional Hilbert spaces. There is no need to limit the physical system to 2-dimensions. Simply choose two appropriate states which form a 2-dimensional subspace in an appropriate basis. The coefficients α and β are the probability amplitudes, and tells with what probability one will find either $|0\rangle$ or $|1\rangle$ when measuring. One finds $|0\rangle, |1\rangle$ with probability $|\alpha|^2$ and $|\beta|^2$ respectively.

The consequence of superposition is that the qubit is not limited to just either 0 or 1, but rather any combination of the two. Superposition together with other QM effects is what is used to create classically impossible algorithms [1, 2]. The combined state of two qubits $|\psi_1\rangle$ and $|\psi_2\rangle$ is given by

$$|\psi_1\rangle \otimes |\psi_2\rangle = (\alpha_1 |0\rangle + \beta_1 |1\rangle) \otimes (\alpha_2 |0\rangle + \beta_2 |1\rangle) \quad (19)$$

the tensor product is often omitted and one would write $|\psi_1\rangle \otimes |\psi_2\rangle = |\psi_1\rangle |\psi_2\rangle = |\psi_1, \psi_2\rangle$. So n qubits is often written as $|x\rangle^{\otimes n}$, $x \in \{0, 1\}$ which is spanned by kets on the form $|x_0 x_1 x_2 \dots x_n\rangle$, $x_i \in \{0, 1\}$, representing all possible bit strings of length n . However we will limit ourselves to single qubit(qudit) systems in this project.

2.2.2 From classical logic gates to quantum gates

A model used in classical computation is the circuit model, which consists of wires and logic gates. The wires carry the bit information, **0**, **1**, and the gates apply different operations. A common gate is the NOT-gate which flips the content of the bit $\text{NOT}(\mathbf{0}) = \mathbf{1}$ and $\text{NOT}(\mathbf{1}) = \mathbf{0}$. Some other common logic gates are shown in Table 1.

Input		AND	NAND	OR
0	0	0	1	0
0	1	0	1	1
1	0	0	1	1
1	1	1	0	1

Table 1: Some common classical logic gates.

To make logic gates suitable for quantum computation the classical logic gate has to be adjusted. A quantum gate has some desired properties

1. Reversibility
2. Conservation of probability

it will turn out that these properties more or less have the same consequence. The first property is not achieved by the classical gates. For example, given the output of the AND-gate it is impossible to know the state of the bits that went into the gate. An output

of **1** only conveys that the inputs were equal, **1,1** or **0,0**, but not which one. Adding another output that simply is one of the inputs can make the gate reversible, if you get a **0** as output, and know that the first input was **0**, then the second input must be a **1**. This extends to all inputs and outputs. In the quantum mechanical context, a quantum logic gate is a quantum operator acting on the state of the qubit. $A|\psi\rangle = |\psi'\rangle$, so the equivalence of the classical reversibility would be the existence of an inverse operator A^{-1} such that the operation could be reversed $A^{-1}|\psi'\rangle = A^{-1}A|\psi\rangle = \mathbf{1}|\psi\rangle = |\psi\rangle$

Conservation of probability does not in itself have any classical analogue due to it being a quantum property. More concretely it means conservation of the inner product. Given the states $U|\psi\rangle = |\psi'\rangle$ and $U|\varphi\rangle = |\varphi'\rangle$, then the inner product of $|\psi\rangle$ and $|\varphi\rangle$ is the same as for $|\psi'\rangle$ and $|\varphi'\rangle$. Thus one can compute

$$\langle\psi'|\varphi'\rangle = \langle\psi|U^\dagger U|\varphi\rangle = \langle\psi|\varphi\rangle \quad (20)$$

and see that this only holds if $U^\dagger U = UU^\dagger = \mathbf{1}$, which happens to be the condition for U to be a unitary operator. From the unitarity it is possible to also see that $U^\dagger = U^{-1}$. The reversibility is also a consequence of the operators being unitary.

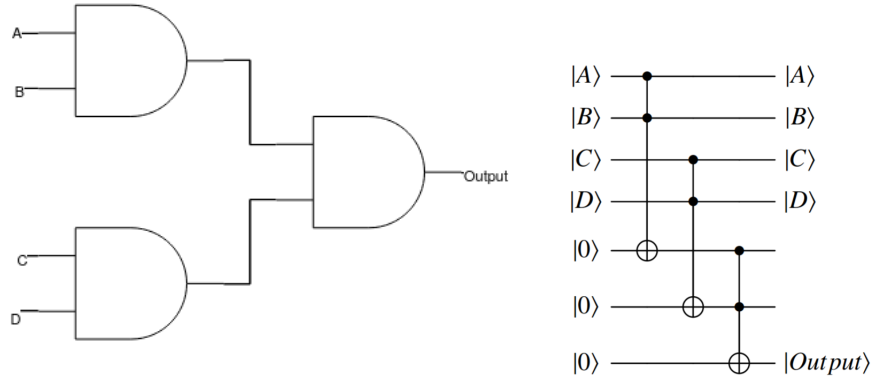


Figure 1: Example of how a classical circuit would be translated into a quantum circuit. The classical AND gates are replaced by the Toffoli gate.

There are many ways to achieve the creation and implementation of quantum gates. Some of them will be discussed in more detail in the Section Geometric Quantum Computation. Let us see in general how a quantum gate could be implemented. To apply a quantum gate to the qubit means to operate on a quantum state with a unitary operator. The time evolution operator discussed priorly is unitary by definition, and it acts on states according to the Schrödinger equation. The time evolution operator seems like a good candidate to be used to implement quantum gates since they share many common properties.

2.2.3 Some important gates and universal computation

Quantum gates are used to apply operations to qubit states. By using gates as standardized building blocks for algorithms we can restrict us to the creation of these gates.

A set of gates is called universal if those gates can be used to approximate any gate to arbitrary precision. Some of the most common gates are the Pauli gates, X, Y, Z , which are equivalent to the Pauli matrices, $\sigma_x, \sigma_y, \sigma_z$. They are defined as

$$X \doteq \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, Y \doteq \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, Z \doteq \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (21)$$

It is more common to use only X and Z since $Y = i X Z$. The dotted equal sign, \doteq , means *is represented by*. The X -gate is somewhat analogous to the classical NOT-gate, with the effect that it swaps the input qubits,

$$\begin{cases} X |0\rangle = |1\rangle \\ X |1\rangle = |0\rangle \end{cases}.$$

The Z -gate introduces a phase-flip between the inputs, it does nothing to $|0\rangle$ and reverses the sign of $|1\rangle$,

$$\begin{cases} Z |0\rangle = |0\rangle \\ Z |1\rangle = -|1\rangle \end{cases}.$$

Superposition is one of the key aspects that make QC able to surpass classical algorithms. The Hadamard gate, H , creates a superimposed quantum state, and is widely used all over QC. The gate is defined as

$$H \doteq \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 1 \\ 1 & -1 \end{pmatrix}. \quad (22)$$

Another noteworthy property is that H is its own inverse, $HH = 1$, so it can be used both to open and close superposition. More explicitly, the gate effect is,

$$\begin{cases} H |0\rangle = \frac{1}{\sqrt{2}}(|0\rangle + |1\rangle) \\ H |1\rangle = \frac{1}{\sqrt{2}}(|0\rangle - |1\rangle) \end{cases} \quad \text{and} \quad \begin{cases} H \frac{1}{\sqrt{2}}(|0\rangle + |1\rangle) = |0\rangle \\ H \frac{1}{\sqrt{2}}(|0\rangle - |1\rangle) = |1\rangle \end{cases}. \quad (23)$$

Many well known quantum algorithms make heavy use of this operator [2, 1], making it essential to QC. The concept of the phase flip, Z , can be extended to a phase shift of any arbitrary angle. The gate is denoted by $P(\phi)$ or sometimes simply ϕ . It is defined as

$$P(\phi) \doteq \begin{pmatrix} 1 & 0 \\ 0 & e^{i\phi} \end{pmatrix}. \quad (24)$$

Then Z reduces to the case $P(\pi) = Z$. Some other notable cases of the phase shift are the T and S gates, defined by

$$T = P(\pi/4) \doteq \begin{pmatrix} 1 & 0 \\ 0 & e^{i\pi/4} \end{pmatrix}, S = T^2 = P(\pi/2) \doteq \begin{pmatrix} 1 & 0 \\ 0 & e^{i\pi/2} \end{pmatrix}. \quad (25)$$

All gates so far are single-qubit gates. To achieve full universality it turns out that two-qubit gates are enough. The two-qubit gate most commonly used is the CNOT, which stands for controlled not, the effect is defined as $\text{CNOT} |x, y\rangle = |x, x \oplus y\rangle$, where \oplus denotes addition modulo 2. The gate takes two inputs, the control, and the target. If the control is $|1\rangle$ then a NOT is applied to the target $|y\rangle \mapsto X |y\rangle$, otherwise the gate leaves the inputs unchanged.

2.2.4 Qudit generalizations

Inherently there is nothing special about bits and binary logic, the same goes for qubits. A qudit is a general name for a d -dimensional computational element. The special case of $d = 3$ is called a *qutrit*, $d = 4$ a *ququart* and so on. A larger computational basis could even prove to be better. The information content per unit is higher, N qubits can be reduced down to $\frac{N}{\log_2(d)}$ qudits [17]. The qutrit has a reduction factor of ~ 0.64 , the ququart 0.5, with less returns for higher dimensions. Due to more information per unit, fewer operations are required to build quantum gates, minimizing the loss of accuracy from operations. The drawback is the increased complexity in the schemes for physical implementation which could come with increased errors. There are already many promising results using qudits [11, 12, 13], and many more are discussed in [14]. The interesting question is if the benefits of qudits can outweigh the extra costs associated with more complex structures.

In the following section, let the dimension d be any integer such that $d \geq 3$. Some generalized qudits gates are defined and discussed. The definitions used are the same used in [14]. The effect of NOT is ambiguous when the number of basis states are more than 2. So instead of viewing it as NOT, it can be thought of as a cyclic permutation of the states $|0\rangle, |1\rangle$. The higher dimensional correspondence of X is defined as

$$X_d = \begin{pmatrix} 0 & 0 & \dots & 0 & 1 \\ 1 & 0 & \dots & 0 & 0 \\ 0 & 1 & \dots & 0 & 0 \\ \vdots & \vdots & \ddots & \vdots & \vdots \\ 0 & 0 & \dots & 1 & 0 \end{pmatrix}. \quad (26)$$

The effect works as a permutation of basis states, for $d = 3$, the effect would be $|0\rangle \xrightarrow{X_3} |1\rangle \xrightarrow{X_3} |2\rangle \xrightarrow{X_3} |0\rangle$. The higher dimensional Z -gate Z_d is defined as

$$Z_d = \begin{pmatrix} 1 & 0 & 0 & \dots & 0 \\ 0 & \omega & 0 & \dots & 0 \\ 0 & 0 & \omega^2 & \dots & 0 \\ \vdots & \vdots & \vdots & \ddots & \vdots \\ 0 & 0 & 0 & \dots & \omega^{d-1} \end{pmatrix} \quad (27)$$

with $\omega = e^{2\pi i/d}$ being the d th root of unity. It shifts the different basis states a certain amount from each other. The generalized version of the Hadamard gate for qudits, H_d , is more convenient to define in Dirac notation, the effect is

$$H_d = \frac{1}{\sqrt{d}} \sum_{i=0}^{d-1} \omega^{ij} |i\rangle \langle j|, \quad j \in \{0, 1, 2, \dots, d-1\} \quad (28)$$

where ω again denotes the d th root of unity.

Last let us define the generalized qudit T -gate, now d is restricted to prime numbers, such that T_d is only defined for prime number dimensions [18]. The gate is defined

by

$$T_d = \sum_{k=0}^{d-1} \omega^{\nu_k} |k\rangle \langle k| \quad (29)$$

with ω being the d th root of unity. The exponent ν_k can be recursively defined with $\gamma', z', \varepsilon' \in \mathbb{Z}_d$ restricted by

$$\begin{cases} \nu_0 = 0 \\ \nu_{k+1} = \nu_k + k(2^{-1}\gamma'k + z') + 2^{-1}z' + \varepsilon'. \end{cases} \quad (30)$$

In this report, T_3 with $z' = 1, \gamma' = 2$ and $\varepsilon' = 0$ will be implemented. This gate reads

$$T_3 = \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{2\pi i/9} & 0 \\ 0 & 0 & e^{-2\pi i/9} \end{pmatrix}. \quad (31)$$

Likewise, as for qubits, single qudit universality can be obtained by a combination of some of these gates. The universal set of gates for the qudit is the generalization of the universal set for the qubit, $\{T_d, H_d\}$. Since the T_d -gate is only defined in prime number dimensions universality cannot be obtained in this manner in non-prime dimensions.

2.2.5 Dark states, bright states and dark paths

Pod-like systems are an important type of quantum system and are used frequently in geometric QC. A pod-like system refers to a type of system with most commonly, but not limited to, one excited state coupled to a number of ground states, states can only be transformed into states to which they are coupled, with no additional coupling between the ground states. This section defines some common terms and concepts that will be used in the later sections.

For a given Hamiltonian, H , the eigenvalues correspond to the energies of the system. Let us denote a state with zero energy as $|d\rangle$. It holds then that

$$H |d\rangle = 0 |d\rangle, \quad (32)$$

which is to say that $|d\rangle$ is an eigenstate with eigenvalue 0. We say that $|d\rangle$ is a **dark state** of H . Quantum states associated with this state is called the dark sub-space of the Hamiltonian. All states in the dark sub-space will have 0 energy. Furthermore we say that any state $|b\rangle$ such that

$$\langle d|b\rangle = 0 \quad (33)$$

is called a **bright state**. A pod-like system with n ground states and m excited states will have $n - m$ dark states and m bright states [19]. The Hamiltonian can be rewritten in the *dark-bright* basis (Morris-Shore basis), an equivalent basis spanned by the bright and dark states. Since the dark states are all individually decoupled from the Hamiltonian this new basis can be further split into two subspaces, one containing the bright states and one containing the dark states. This reduction can be useful since it simplifies the dynamics of the system. The dynamics of a system are closely related to

the Hamiltonian of the system. The part of the time evolution affected by the Hamiltonian, for some ordered basis, is $H_{nm}(t) = \langle n | H(t) | m \rangle$. In the adiabatic case if the system starts in the dark subspace, it will remain there through-out the evolution and the dynamical phase will have no contribution to the total time evolution. In a non-adiabatic case the time evolution of the states is not as well behaved and the states will mix. So by constructing a Hamiltonian such that all H_{nm} elements are zero there will be no dynamical contribution. Another way to assure no dynamical phase is to evolve the system along a **dark path**, a dark path is a state $|D(t)\rangle$ such that

$$\langle D(t) | H(t) | D(t) \rangle = 0, \forall t \in [0, T]. \quad (34)$$

this means that the expectation value of the energy will always be 0, and thus all elements of the dynamical phase will be 0. The dark state in itself is a dark path. By explicitly finding a dark path the parameters in the Hamiltonian can be reverse-engineered from the Schrödinger equation, enabling more complex Hamiltonians that can fulfill the NHQC criteria.

2.2.6 Geometric quantum computation

By making use of the geometric phases in periodic quantum systems schemes that perform computations are obtainable. Holonomic Quantum Computation (HQC) makes use of holonomies, which is to what extent parallel transport around some closed-loop fails to preserve the transported data. In the case of Quantum mechanics this means that a state, gains an additional geometric phase alongside the dynamical phase. This concept was first introduced by M.V. Berry [5] and thus the phase is usually referred to as the Berry phase but is more generally called the geometric phase since the phase arises from the geometry rather than the dynamics of the system. The geometric phase exists also in non-adiabatic and non-abelian contexts, which gives room for more complex structures [6, 7, 8]. In the adiabatic setting universal computation can be achieved through holonomies by transporting a set of control parameters along a loop(s) in a suitable space [3].

For an arbitrary periodic state $|\psi(T)\rangle$ after one period T , the Schrödinger equation (16) has a solution on the form

$$|\psi(T)\rangle = U(T, 0) |\psi(0)\rangle \quad (35)$$

where in addition to the ordinary dynamical phase discussed earlier a geometric phase is present. For an ordered time dependent basis $|1(t)\rangle, |2(t)\rangle, \dots, |N(t)\rangle$ a quantum state can be expanded into

$$|\psi\rangle = \sum_{n=1}^N c_{:n} |n\rangle \quad (36)$$

taking the time derivative gives

$$|\dot{\psi}\rangle = \sum_{n=1}^N \left[\dot{c}_{:n} |n\rangle + c_{:n} |\dot{n}\rangle \right]. \quad (37)$$

By the Schrödinger equation $|\dot{\psi}\rangle$ is equivalent to

$$|\dot{\psi}\rangle = -i H |\psi\rangle = -i \sum_{n=1}^N c_n H |n\rangle \quad (38)$$

equating (37) and (38) and taking the inner product with $\langle m|$ one obtains the expression

$$-i c_{mn} \langle m| H |n\rangle = \dot{c}_{mn} + \sum_{n=1}^N c_{mn} \langle m|\dot{n}\rangle \quad (39)$$

re-organizing terms we get

$$\dot{c}_{mn} = -i \left[c_{mn} \langle m| H |n\rangle - \sum_{n=1}^N i c_{mn} \langle m|\dot{n}\rangle \right]. \quad (40)$$

In matrix form this can be written as

$$\dot{\mathbf{C}} = i [\mathbf{A} - \mathbf{H}] \mathbf{C}, \quad (41)$$

where $\mathbf{A}_{mn}(t) = i \langle m(t)| \frac{\partial}{\partial t} |n(t)\rangle$ and $\mathbf{H}_{mn}(t) = \langle m(t)| \mathbf{H}(t) |n(t)\rangle$. The formal solution is

$$C = U(T, 0) = \mathcal{T} \exp \left(i \int_0^T \underbrace{\mathbf{A}(t)}_{\text{Geometric}} - \underbrace{\mathbf{H}(t)}_{\text{Dynamic}} dt \right) \quad (42)$$

If a system is periodic and a scheme can be constructed such that the dynamical phase will have no contribution to the time evolution during a full period, the period can also be described as a loop, C , in parameter space. So for a period the the operator can be written as

$$U(T, 0) = \mathcal{T} \exp \left(i \int_0^T \mathbf{A}(t) dt \right) = \mathcal{P} \exp \left(i \oint_C \mathbf{A}(s) ds \right) = U(C), \quad (43)$$

where \mathcal{P} denotes path-ordering. This is the Wilczek-Zee holonomy which can be used to achieve geometric quantum computation.

In the adiabatic case it is enough for the system to start in a dark sub-space since the states will evolve without mixing. Thus the energy associated with the Hamiltonian will be zero through-out the whole time evolution. Therefore it is possible to construct unitaries by pure geometric means by starting in the correct state [20, 21].

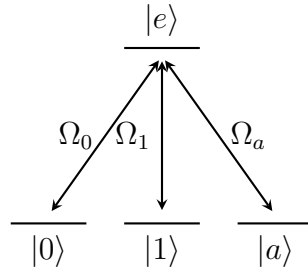


Figure 2: System described by the Hamiltonian in Equation 44.

A well known result that makes use of the Wilczek-Zee holonomy to obtain universal computation presented by L.M. Duan in [21]. By using a tripod system defined by the Hamiltonian

$$H = |e\rangle [\Omega_0 \langle 0| + \Omega_1 \langle 1| + \Omega_a \langle a|] + \text{h.c.}, \quad (44)$$

the system is shown in Figure 2. By manipulating the parameters and the coupling strengths this Hamiltonian can be used to create a set of universal gates; $U_1 = e^{i\phi_1|1\rangle\langle 1|} = \begin{pmatrix} 1 & 0 \\ 0 & e^{i\phi_1} \end{pmatrix}$, $U_2 = e^{i\phi_2\sigma_y} = R_y(\phi_2)$, $U_3 = e^{i\phi_3|11\rangle\langle 11|}$. Let us calculate U_1 , which is a phase gate, to see how this is obtained. First the parameters are chosen as $\Omega_0 = 0$, $\Omega_1 = -\Omega \sin \frac{\theta}{2} e^{i\varphi}$, $\Omega_a = \Omega \cos \frac{\theta}{2}$. The Hamiltonian in this case is

$$H = |e\rangle \left[-\Omega \sin \frac{\theta}{2} e^{i\varphi} \langle 1| + \Omega \cos \frac{\theta}{2} \langle a| \right] + \text{h.c.}, \quad (45)$$

where + h.c is read as Hermitian conjugate, thus $A + \text{h.c} = A + A^\dagger$. Find a dark state $|d\rangle$ such that $H|d\rangle = 0$. In this state the energy will be zero and by evolving the system adiabatically the energy will remain zero. The dark state is $|d\rangle = \cos \frac{\theta}{2} |1\rangle + \sin \frac{\theta}{2} e^{i\varphi} |a\rangle$. By starting in the state $\theta = 0$, then $|d\rangle = |1\rangle$ and evolving the new control parameters in a closed loop in parameter space we find that the holonomy is independent of the path and $U_1 = e^{i\phi_1|1\rangle\langle 1|}$ where ϕ_1 is the enclosed solid angle obtained by the vector (θ, φ) during the evolution, $\phi_1 = \oint \sin \theta d\theta d\varphi$. The other single-qubit gate U_2 is obtained in a similar way and U_3 requires some extra work by using detuning to create a two-qubit gate.

Due to the long runtimes associated with adiabatic evolution the state is prone to errors and decoherence from outside factors [4]. As a consequence, in the non-adiabatic case (NHQC), it would be possible to construct faster and more robust gates. Since the states are mixing some other way to control U must be engineered. By removing the dependence on the dynamics of the system, the period time, and the energies in the time evolution, that is by some other mean ensure that

$$\langle m(t) | H(t) | n(t) \rangle = 0, \quad \forall m, n, t \in [0, T] \quad (46)$$

holds. This means all terms in the dynamical phase in (42) are zero. This is possible given that there exist a two-pulse non commuting loop, $C = C_1 \cup C_2$, [4, 22], then the unitary will be reduced to

$$U(T, 0) = \sum_{n,m=1}^N U_{nm} |n(0)\rangle \langle m(0)| \quad (47)$$

and can be used to perform any arbitrary quantum gate on the computational space.

The idea of NHQC can be extended by instead of explicitly constructing a Hamiltonian with non-dynamical interaction, one can construct explicit dark paths through parameter space that satisfy the NHQC criteria. In [15] a qubit scheme is constructed and experimentally implemented in a trapped $^{171}\text{Yb}^+$ ion, using a dark path. They conclude that arbitrary NHQC gates can be constructed with better robustness. We will build upon the scheme presented and create an explicit qutrit example, then speculate how it can be generalized to higher dimensions.

3 Results

3.1 The qutrit

3.1.1 Construction of the system

In [15] a qubit is implemented by modifying the lambda system, adding an additional level, and an auxiliary state, effectively turning it into a tripod. The generalization to implement a qutrit is not trivial but it is clear that one extra ground state is required to make the computational basis larger. Another thing that is needed is to limit the number of dark states to one. The number of dark states is the difference between the number of excited states and the number of ground states [19] (here the auxiliary state is **not** counted as a ground state). Therefore another excited state must be added as well. How these states are coupled is the hard part. We will see that the system given by the Hamiltonian

$$H = \sum_{j=1}^2 \sum_{i=j}^3 \omega_{ij} |i\rangle \langle e_j| + \frac{\Omega_a(t)}{2} |a\rangle \langle e_2| + \text{h.c} \quad (48)$$

which is shown in Figure 3, will work. It consists of two excited states, $|e_1\rangle, |e_2\rangle$, an auxiliary state $|a\rangle$ and the ground states $|1\rangle, |2\rangle, |3\rangle$ which makes up the computational basis of the qutrit. By changing the basis of the Hamiltonian, it can be rewritten as

$$H_d = \sum_{j=1}^2 \frac{\Omega_j(t)}{2} e^{-i\phi_j} |b_j\rangle \langle e_j| + \frac{\Omega_a(t)}{2} |a\rangle \langle e_2| + \text{h.c} \quad (49)$$

by a Morris-Shore transformation [23]. This new basis is the dark-bright basis. To find this basis find a dark state to the original Hamiltonian, an eigenstate with eigenvalue 0, $H|d\rangle = 0$, then find two bright states, orthogonal to both $|d\rangle$ and each other. This basis can be parametrized with 4 angles, $\theta, \varphi, \chi, \xi$, the dark state is on the form $|d\rangle = \cos \theta |1\rangle + e^{i\chi} \sin \theta \cos \varphi |2\rangle + e^{i\xi} \sin \theta \sin \varphi |3\rangle$, additionally the two bright states are $|b_1\rangle = N_1 (-e^{-i\chi} \sin \theta \cos \varphi |1\rangle + \cos \theta |2\rangle)$ and $|b_2\rangle = N_2 (\cos \theta |1\rangle + e^{i\chi} \sin \theta \cos \varphi |2\rangle + \Lambda |3\rangle)$, where N_1, N_2 are normalization factors and Λ can be chosen such that $\langle d|b_2\rangle = 0$.

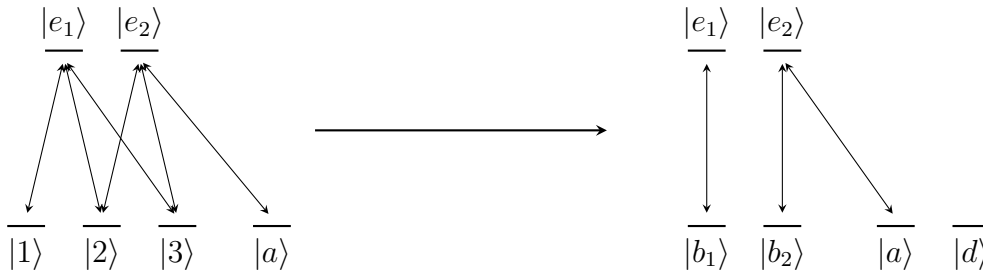


Figure 3: The system given by the Hamiltonian shown in Equation 48 (left) and the transformed system from Equation 49 (right)

Explicitly the states in the dark-bright basis are given by

$$\begin{aligned}
|d\rangle &= \cos \theta |1\rangle + e^{i\chi} \sin \theta \cos \varphi |2\rangle + e^{i\xi} \sin \theta \sin \varphi |3\rangle \\
|b_1\rangle &= \frac{1}{\sqrt{1 - \sin^2 \theta \sin^2 \varphi}} (-e^{-i\chi} \sin \theta \cos \varphi |1\rangle + \cos \theta |2\rangle) \\
|b_2\rangle &= \frac{1}{\sqrt{1 - \sin^2 \theta \sin^2 \varphi}} \left(\frac{1}{2} \sin 2\theta \sin \varphi |1\rangle + \frac{e^{i\chi}}{2} \sin^2 \theta \sin 2\varphi |2\rangle + e^{i\xi} (\sin^2 \theta \sin^2 \varphi - 1) |3\rangle \right)
\end{aligned} \tag{50}$$

The parameters ω_{ij} in the original basis can be determined by replacing the states in H_d by their form in the $\{|1\rangle, |2\rangle, |3\rangle\}$ basis.

Now let's introduce two dark paths, $|D_1(t)\rangle, |D_2(t)\rangle$. Along these paths the average energy is always zero $\langle D_i(t) | H_d | D_i(t) \rangle = 0, i = 1, 2$. Thus no dynamical phase is accumulated during the time evolution and therefore it follows the conditions required for NHQC.

The following two states satisfy the dark path condition and can be parametrized by two angles $u(t), v(t)$.

$$\begin{aligned}
|D_1(t)\rangle &= \cos u e^{-i\phi_1} |b_1\rangle + i \sin u |e_1\rangle \\
|D_2(t)\rangle &= \cos u \cos v e^{-i\phi_2} |b_2\rangle - i \sin u |e_2\rangle - \cos u \sin v |a\rangle
\end{aligned} \tag{51}$$

it can easily be verified that $\langle D_i(t) | H_d | D_j(t) \rangle = 0, i, j = 1, 2$. The angles can be chosen with the constraint $|D_i(0)\rangle \langle D_i(0)| = |D_i(T)\rangle \langle D_i(T)|, i = 1, 2$. This can be achieved by choosing $u(0) = u(T) = v(0) = v(T) = 0$. A valid choice is $u(t) = \frac{\pi}{2} \sin^2 \frac{\pi t}{T}$ and $v(t) = \eta [1 - \cos u(t)]$, as for the qubit case [15]. η represents the coupling strength to the auxiliary state $|a\rangle$ and the system reverts into a bi-tri-pod structure. Unless mentioned otherwise $\eta = 4.0$. This is the optimal choice for the qubit and shall be used here as well [15]. Each dark path starts in the respective bright state and travels along a curve and then returns to the bright state. Using the Schrödinger equation one can relate the dark path to the Hamiltonian,

$$i \frac{\partial}{\partial t} |D_i(t)\rangle = H_d |D_i(t)\rangle, \tag{52}$$

and thusly one can reverse engineer the time dependent parameters $\Omega_i(t)$ by matching the factors in front of each state. A calculation yields

$$\begin{aligned}
\Omega_1(t) &= -2\dot{u} \\
\Omega_2(t) &= 2(\dot{v} \cot u \sin v + \dot{u} \cos v) \\
\Omega_a(t) &= 2(\dot{v} \cot u \cos v - \dot{u} \sin v).
\end{aligned} \tag{53}$$

To construct a quantum gate, we make use of the method of multi-pulse single-loops [22], the relevant part of the time evolution operator is

$$\begin{aligned}
U_1 &= |d\rangle \langle d| - i |e_1\rangle \langle b_1| - i |e_2\rangle \langle b_2|, \phi_1 = \phi_2 = 0 \\
U_2 &= |d\rangle \langle d| + i e^{i\gamma_1} |b_1\rangle \langle e_1| + i e^{i\gamma_2} |b_2\rangle \langle e_2|, \phi_1 = -\gamma_1, \phi_2 = -\gamma_2
\end{aligned} \tag{54}$$

so the operator for one full loop is

$$U = U_2 U_1 = |d\rangle \langle d| + e^{i\gamma_1} |b_1\rangle \langle b_1| + e^{i\gamma_2} |b_2\rangle \langle b_2|. \tag{55}$$

This transformation can be parametrized by 6 real parameters, $\chi, \xi, \theta, \varphi, \gamma_1, \gamma_2$, however it is not enough to construct all gates. For example X_3 requires 2 loops. This is since one loop does not cover all degrees of freedom. The reason for this is elaborated on later in the Generalization section. The full gate is given by repeating U with another set of parameters. Thus full gate \mathbb{U} is given by

$$\mathbb{U} = U(\chi', \xi', \theta', \varphi', \gamma'_1, \gamma'_2)U(\chi, \xi, \theta, \varphi, \gamma_1, \gamma_2) \quad (56)$$

Now the problem is only a matter of finding all parameters to replicate the desired gate. This is a non-trivial problem since it includes solving a system of 9 non-linear algebraic equations containing 12 variables for a gate requiring two loops. Some gates that only require a single loop can with some work be found analytically but the problem quickly becomes complex.

3.1.2 Implementation

The implementation is quite straightforward. For a given set of parameters, say that we want to see how the initial state $|\psi(0)\rangle$ evolves with time. Let $|\psi(0)\rangle = (a, b, c)^T$ in the computational basis $\{|1\rangle, |2\rangle, |3\rangle\}$. Then the following equation can be formulated, for some coefficients f_i ,

$$\begin{pmatrix} a \\ b \\ c \end{pmatrix} = f_0 |d\rangle + f_1 |b_1\rangle + f_2 |b_2\rangle \quad (57)$$

since at time $t = 0$ the dark paths are in the corresponding bright state. Now expand the dark state and bright state in terms of the original basis. Doing this one can obtain the relation

$$\begin{pmatrix} f_0 \\ f_1 \\ f_2 \end{pmatrix} = \begin{pmatrix} c_1 & -N_1 c_2 & N_2 c_1 \\ c_2 & N_1 c_1 & N_2 c_2 \\ c_3 & 0 & N_2 \Lambda \end{pmatrix}^{-1} \begin{pmatrix} a \\ b \\ c \end{pmatrix} \quad (58)$$

So by choosing an initial state, the coefficients can be obtained. Then the system can be solved for any time $t \in T$ using Equation ??

$$|\psi(t)\rangle = f_0 |d\rangle + f_1 |D_1(t)\rangle + f_2 |D_2(t)\rangle, t \in [0, T]. \quad (59)$$

To see the effect of the gates, the probability of finding the system in the standard computational basis states can be plotted against time to see how the gate transform the state from initial to final state, shown in Figure 4 and 5 for the X_3 and H_3 gates respectively .

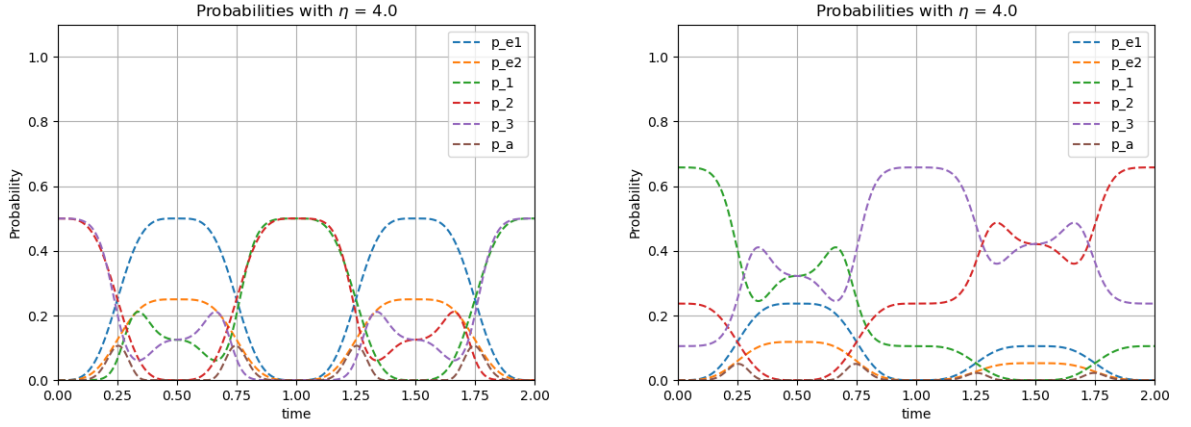


Figure 4: The effect of the X_3 -gate on the initial states $\frac{1}{\sqrt{2}}[0, 1, 1]$ (left) and $\frac{1}{\sqrt{38}}[5, 3, 2]$ (right). Note that since the plot shows the probabilities, phases cannot be seen in the plot.

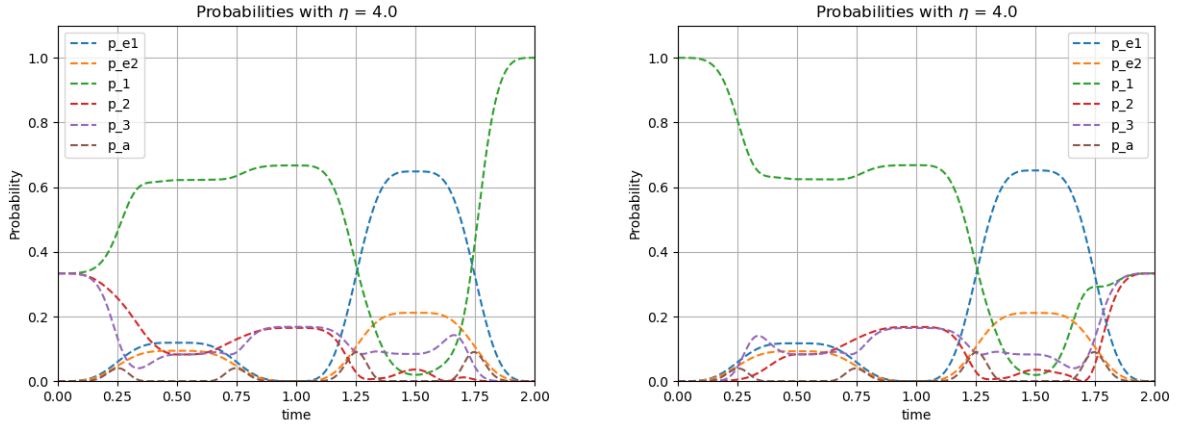


Figure 5: The effect of the H_3 -gate on the initial states $\frac{1}{\sqrt{3}}[1, 1, 1]$ (left) and $[1, 0, 0]$ (right). Note that since the plot shows the probabilities, phases cannot be seen in the plot.

3.1.3 Important gates, universality and parameter selection

By specifying the parameters of U the qutrit gates from the Background section can be constructed. A selection of important gates are obtained by the following;

$$\begin{aligned}
X_3 &= U(0, 0, \frac{\pi}{4}, \frac{\pi}{2}, 0, \pi) \times U(0, 0, \frac{\pi}{2}, \frac{\pi}{4}, 0, \pi) = \begin{pmatrix} 0 & 0 & 1 \\ 1 & 0 & 0 \\ 0 & 1 & 0 \end{pmatrix} \\
Z_3 &= U(0, 0, 0, 0, \frac{2\pi}{3}, \frac{4\pi}{3}) = \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{\frac{2\pi i}{3}} & 0 \\ 0 & 0 & e^{\frac{4\pi i}{3}} \end{pmatrix} \\
T_3 &= U(0, 0, 0, 0, \frac{2\pi}{9}, \frac{-2\pi}{9}) = \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{\frac{2\pi i}{9}} & 0 \\ 0 & 0 & e^{\frac{-2\pi i}{9}} \end{pmatrix} \\
H_3 &= U(6.41 \cdot 10^{-4}, 6.56 \cdot 10^{-4}, 0.48, 0.79, 1.58, 1.56) \\
&\times U(9.81 \cdot 10^{-3}, 0.00, 1.187, 2.15, 0.00, 1.57) \\
&\approx \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 1 & 1 \\ 1 & e^{\frac{2\pi i}{3}} & e^{\frac{4\pi i}{3}} \\ 1 & e^{\frac{4\pi i}{3}} & e^{\frac{2\pi i}{3}} \end{pmatrix}.
\end{aligned} \tag{60}$$

Note that the choice of parameters is not unique and there are multiple ways to create the same unitary. These gates are enough to achieve single-qutrit universality as discussed in Section Theoretical Background. To achieve full universality an additional two-qutrit gate is needed. A common choice for qubits is the CNOT gate [14].

The parameters for any given gate can be specified by solving the non-linear system of equations obtained from by setting the matrix representation of the gate equal to Equation 55 or 56. It can be found that some gates such as the X_3 -gate can not be replicated by only one loop but will require at least 2 loops. The physical effect of this is that during the first cycle a pair states are swapped, and during the second cycle another pair is swapped. This can be seen in Figure 4. For the diagonal gates the analytical solution can be easily found by fixing $\theta = \varphi = \chi = \xi = 0$. The basis states reduce to $|d\rangle = |1\rangle$, $|b_1\rangle = |2\rangle$, $|b_2\rangle = -|3\rangle$. Then all diagonal unitaries can be specified by γ_1 and γ_2 up to a phase factor,

$$U(0, 0, 0, 0, \gamma_1, \gamma_2) = \begin{pmatrix} 1 & 0 & 0 \\ 0 & e^{i\gamma_1} & 0 \\ 0 & 0 & e^{i\gamma_2} \end{pmatrix}. \tag{61}$$

Analytical solutions are not always easy to find. The H_3 -gate gives rise to a system of equations that is non-trivial to solve. In that case an approximative gate \tilde{U} can be found by numerical optimization of the expression $\min ||U - \tilde{U}||_F$, where $||\cdot||_F$ is the Frobenius matrix norm given by $||A||_F = \sqrt{\text{Tr}[A^\dagger A]}$. Since $\tilde{U} \approx U$ the approximated gate will have an effect close to that of the exact gate.

To assess the robustness of the gate the fidelity metric, F , is used, the metric measures how close two quantum states are to each other. Here we are only considering pure states, then the fidelity is: $F(|\psi\rangle, |\varphi\rangle) = |\langle\psi|\varphi\rangle|$. The fidelity is averaged by

sampling initial states and letting them evolve with time by numerically solving the Schrödinger Equation using the SciPy implementation of backwards differentiation [24]. Then introducing small shifts, $\Omega \mapsto \Omega(1 + \delta)$, for all Ω_i pulses in Equation 49 and then comparing to the exact solution obtained by multiplying the gate with the initial state. The calculated fidelities are shown in Figure 6. We see that the coupling with the auxiliary state is more resilient to errors in the parameters than the original NHQC scheme. This is very similar to the results in 2 dimensions from [15]. Which suggests an improvement of robustness even in higher dimensions.

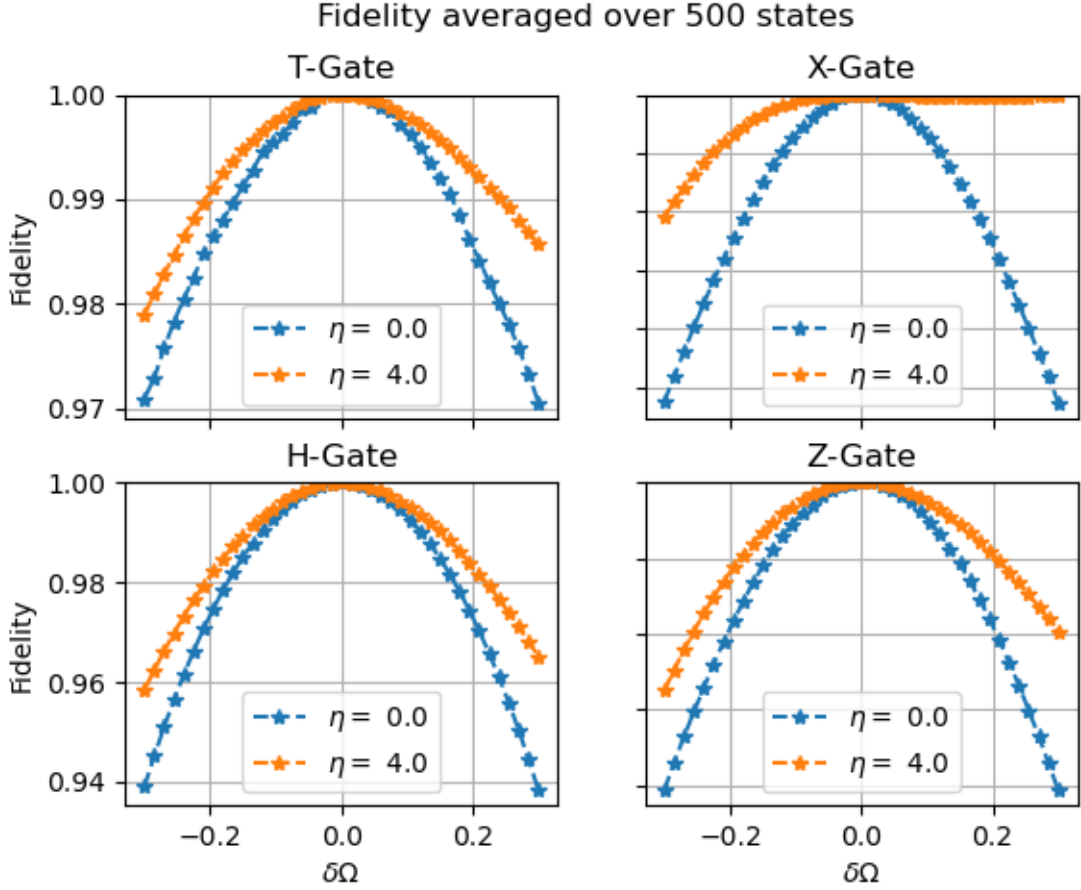


Figure 6: Robustness test, average fidelity of the T_3 , X_3 , Z_3 , and H_3 gates. Average calculated by sampling over 500 randomized initial states with a perturbation of the Ω -pulses, $\Omega \mapsto \Omega(1 + \delta)$.

3.2 Generalization

To generalize the scheme for qudits with arbitrary dimension n , the same idea from the qutrit case can be extended, n ground states, $m = n - 1$ excited states and 1 auxiliary state, therefore the dimension of the Hilbert space is $n + m + 1 = n + 1 - 1 = 2n$. Once again the couplings are not trivial and will be somewhat intricate, but the couplings given by the Hamiltonian in Equation 62, will do the trick.

$$H = \sum_{j=1}^m \sum_{i=j}^n \omega_{ij} |i\rangle \langle e_j| + \frac{\Omega_a(t)}{2} |a\rangle \langle e_m| + \text{h.c} \quad (62)$$

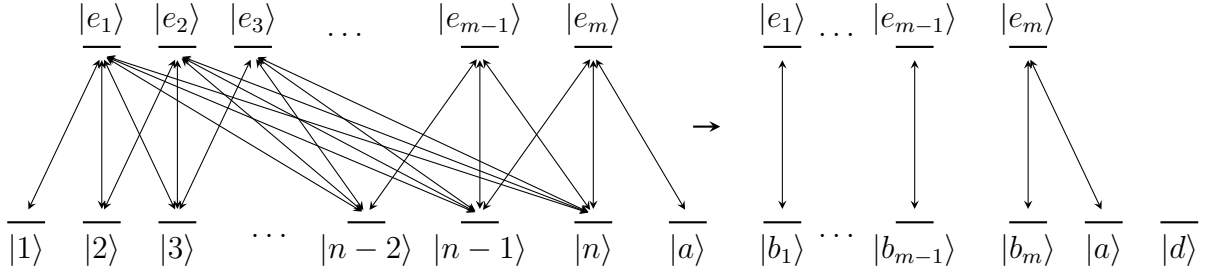


Figure 7: The system given by the Hamiltonian shown in Equation 62 (left) and the transformed system from Equation 69 (right)

The general system for a qudit is given by m excited states $|e_i\rangle$, $i = 0, 1, \dots, m$ and n ground states labeled $|i\rangle$, $i = 1, 2, \dots, n$ and one auxiliary state $|a\rangle$. The number of states always follows that $n - m = 1$ to limit the number of dark states to one [19]. The transitions occur only between excited states and ground states. The couplings follow a pattern, the first excited state e_1 is connected to all ground states, e_2 is connected to all ground states except $|1\rangle$, in general, the excited state $|e_i\rangle$ is connected to the $(n - i)$ th ground states with the largest index. Unless $i = m$, the excited state $|e_m\rangle$ is connected to the two highest labeled ground states and the auxiliary state $|a\rangle$. See Equation 62 and Figure 7 for a clarification.

From the standard basis $\{e_1, e_2, \dots, e_m, 1, 2, \dots, n, a\}$, it is possible to change to a new basis, the dark-bright basis given by $\{e_1, e_2, \dots, e_m, b_1, b_2, \dots, b_m, a, d\}$. Assume there exists a dark state on the form

$$|d\rangle = c_1 |1\rangle + c_2 |2\rangle + c_3 |3\rangle \cdots + c_n |n\rangle, \quad |c_1|^2 + |c_2|^2 + \cdots + |c_n|^2 = 1 \quad (63)$$

from this dark state one could define $n - 1 = m$ bright states. Starting from $|b_1\rangle$, constructing it such that it will be orthogonal to the dark state and all other bright states,

$$|b_1\rangle = N_1 (-c_2^* |1\rangle + c_1^* |2\rangle) \quad (64)$$

with N_1 being a normalization factor. Then choose additional bright states on the form

$$\begin{aligned}
|b_2\rangle &= N_2 (c_1 |1\rangle + c_2 |2\rangle + \Lambda_3 |3\rangle) \\
|b_3\rangle &= N_3 (c_1 |1\rangle + c_2 |2\rangle + c_3 |3\rangle + \Lambda_4 |4\rangle) \\
&\vdots \\
|b_{m-1}\rangle &= N_{m-1} (c_1 |1\rangle + c_2 |2\rangle + c_3 |3\rangle + \cdots + c_{m-1} |m-1\rangle + \Lambda_m |m\rangle) \\
|b_m\rangle &= N_m (c_1 |1\rangle + c_2 |2\rangle + c_3 |3\rangle + \cdots + c_m |m\rangle + \Lambda_{m+1} |m+1\rangle).
\end{aligned} \tag{65}$$

By construction, $|b_1\rangle$ is orthogonal to the dark state all other bright states. for $k \geq 2$ $|b_k\rangle$ is $(k+1)$ -dimensional, so it consists of $k+1$ basis-kets, where the coefficient in front of the last ket, Λ_{k+1} is chosen such that $|b_k\rangle$ is orthogonal the dark state, $|d\rangle$. This will in turn make $|b_k\rangle$ orthonormal to any $|b_{>k}\rangle$ as they have the same states and coefficients as $|d\rangle$ for all the states involved in the inner product $\langle b_{>k} | b_k \rangle = \langle d | b_k \rangle$. Therefor by chosing Λ_k such that these inner product are zero the construction will results in orthonormal states, $\langle b_i | b_j \rangle = \delta_{ij}$. Explicitly for the state $|b_k\rangle, k \geq 2$, the coefficient Λ_{k+1} will be given by

$$\Lambda_{k+1} = -\frac{1}{c_{k+1}^*} \sum_{l=1}^k |c_l|^2 \tag{66}$$

and the normalization N_k is given by

$$N_k = \left(\sum_{l=1}^k |c_l|^2 + |\Lambda_{k+1}|^2 \right)^{-1/2} = \left(\sum_{l=1}^k |c_l|^2 + \left| -\frac{1}{c_{k+1}^*} \sum_{l=1}^k |c_l|^2 \right|^2 \right)^{-1/2} \tag{67}$$

The coefficients c_i can be parametrized by the euclidean components of the radius of the unit- n -sphere with an added phase factor.

$$\begin{aligned}
c_1 &= \cos(\varphi_1) \\
c_2 &= e^{i\theta_1} \sin(\varphi_1) \cos(\varphi_2) \\
c_3 &= e^{i\theta_2} \sin(\varphi_1) \sin(\varphi_2) \cos(\varphi_3) \\
&\vdots \\
c_{n-1} &= e^{i\theta_{n-1}} \sin(\varphi_1) \dots \sin(\varphi_{n-2}) \cos(\varphi_{n-1}) \\
c_n &= e^{i\theta_n} \sin(\varphi_1) \dots \sin(\varphi_{n-2}) \sin(\varphi_{n-1})
\end{aligned} \tag{68}$$

c_1 does not need a phase factor since the overall phase of a state is non-measurable and can be chosen such that the first phase factor will cancel out.

In the dark-bright basis, the Hamiltonian can be written as

$$H_d = \sum_{i=1}^m \frac{\Omega_i(t)}{2} e^{-i\phi_i} |b_i\rangle \langle e_i| + \frac{\Omega_a(t)}{2} |a\rangle \langle e_m| + \text{h.c} \tag{69}$$

with Ω_i being real-valued time dependent parameters and the ϕ_i time independent phase factors. The explicit transformation from $H \mapsto H_d$ is not that important but could be obtained by expanding the bright state kets into the standard basis from (65)

With this Hamiltonian m independent dark paths can be constructed, and must satisfy $\langle D_i(t) | H_d | D_i(t) \rangle = 0, i = 1, 2, \dots, m$ and $\langle D_i(t) | D_j(t) \rangle = \delta_{ij}$. The dark paths can be parametrized by two functions $u(t), v(t)$ that satisfy the conditions $u(0) = v(0) = u(T) = v(T) = 0$, will have the form

$$\begin{aligned} |D_i(t)\rangle &= \cos u e^{-i\phi_i} |b_i\rangle + i \sin u |e_i\rangle, \quad i = 1, 2, \dots, m-1 \\ |D_m(t)\rangle &= \cos u \cos v e^{-i\phi_n} |b_m\rangle - i \sin u |e_m\rangle - \cos u \sin v |a\rangle \end{aligned} \quad (70)$$

The dark paths start in the bright state and travel along a curve where the expectation value of the energy is constantly 0 and can thus be used non-adiabatically.

By using these states one can reverse engineer the Hamiltonian using the Schrödinger equation to determine Ω_i and Ω_a since

$$i \frac{\partial}{\partial t} |D_i(t)\rangle = H_d |D_i(t)\rangle, \quad i = 1, 2, \dots, m \quad (71)$$

a calculation yields

$$\begin{aligned} \Omega_1(t) &= -2\dot{u} \\ \Omega_2(t) &= -2\dot{u} \\ &\vdots \\ \Omega_{m-1}(t) &= -2\dot{u} \\ \Omega_m(t) &= 2(\dot{v} \cot u \sin v + \dot{u} \cos v) \\ \Omega_a(t) &= 2(\dot{v} \cot u \cos v - \dot{u} \sin v) \end{aligned} \quad (72)$$

The time evolution is split into k loops, each loop with two pulses, $0 \rightarrow T/2$ and $T/2 \rightarrow T$. The relevant part of the time evolution operator for one loop is

$$U_1(T/2, 0) = |d\rangle \langle d| - i \sum_{i=1}^m |e_i\rangle \langle b_i|, \quad \phi_i = 0, i = 1, 2, \dots, n \quad (73)$$

and

$$U_2(T, T/2) = |d\rangle \langle d| + i \sum_{i=1}^m e^{i\gamma_i} |b_i\rangle \langle e_i|, \quad \phi_i = 0, i = 1, 2, \dots, n \quad (74)$$

The full operator for one loop is then given by

$$U(T, 0) = U_2 U_1 = |d\rangle \langle d| + \sum_{i=1}^m e^{i\gamma_i} |b_i\rangle \langle b_i|. \quad (75)$$

It is clear that U is unitary in the subspace $\{d, b_1, b_2, \dots, b_n\}$. The unitary is parametrized by $3m = 3(n-1)$ parameters for $n \geq 2$,

$$U(\varphi_1, \dots, \varphi_m, \theta_1, \dots, \theta_m, \gamma_1, \dots, \gamma_m). \quad (76)$$

Applying the unitary with different parameters in sequence up to k times is enough to create any desirable gate $\mathbb{U} = U^k$. The unitary is controlled by $3(n-1)$ parameters while n dimensional qudit has more degrees of freedom than covered with a single loop.

The qudit state space is equivalent to $SU(n)$, which has dimensionality $\dim(SU(n)) = n^2 - 1$. To cover all degrees of freedom k must satisfy $3(n-1)k \geq n^2 - 1$. which means $k \geq \frac{n+1}{3}$. Thus the number of loops needed to create any unitary scales linearly at worst since some gates can be created with fewer loops, see Table 2.

In the case of equality $k = \frac{n+1}{3}$, when $n = 3j + 2$, $j \in \mathbb{N}$, the fewest amount of loops per dimension is achieved and could potentially be more efficient carriers of information since the same number of loops must be carried out while higher dimension has higher information capacity.

n	$3(n-1)$	$n^2 - 1$	k	'='
2	3	3	1	✓
3	6	8	2	
4	9	15	2	
5	12	24	2	✓
6	15	35	3	
7	18	48	3	
8	21	63	3	✓
9	24	80	4	
10	27	99	4	
11	30	120	4	✓
12	33	143	5	

Table 2: Table for some dimensions.

The scaling of the qudit gate is linear at worst both in the number of loops and the number of parameters needed for control. In fact, any diagonal gate only requires one loop, by setting $\varphi_1 = \dots \varphi_m = \theta_1 = \dots = \theta_m = 0$ the unitary reduces to the form

$$U(0, \dots, 0, \gamma_1, \dots, \gamma_m) = |1\rangle \langle 1| + \sum_{k=2}^n e^{i\gamma_k} |k\rangle \langle k|. \quad (77)$$

The effect of this variable choice makes $c_1 = 1, c_{i \neq 1} = 0$. By (65) the only thing not clear is how this affects the Λ coefficients. For the state $|b_k\rangle$, $k \geq 3$, fixing $c_1 = 1$ and let all other states approach 0.

$$\begin{aligned}
\lim_{c_{i \neq 1} \rightarrow 0} |b_k\rangle &= \lim_{c_{i \neq 1} \rightarrow 0} \frac{1}{N_k} \Lambda_{k+1} |k+1\rangle \\
&= \lim_{c_{i \neq 1} \rightarrow 0} \left(1 + \left| -\frac{1}{c_{k+1}^*} \right|^2 \right)^{-1/2} \left(-\frac{1}{c_{k+1}^*} \right) |k+1\rangle \\
&= \lim_{c_{i \neq 1} \rightarrow 0} \left| -\frac{1}{c_{k+1}^*} \right|^{-1} \left(-\frac{1}{c_{k+1}^*} \right) |k+1\rangle \\
&= -|k+1\rangle
\end{aligned} \quad (78)$$

Therefor $|d\rangle = |1\rangle, |b_1\rangle = |2\rangle, |b_2\rangle = -|3\rangle, |b_3\rangle = -|4\rangle, \dots, |b_m\rangle = -|m+1\rangle$, and thus the unitary will take the form (77) so any diagonal unitary can be created up to any unimportant phase.

Thusly, any phase flip gate, such as Z_n can be constructed in any dimension using only one loop. Since the gates essential for single-qudit universality is $\{T_n, H_n\}$, where one of the gates can be obtained by a single loop. Thus only the Hadamard gate will be the most costly part of the scheme as it requires at most $k = \frac{n+1}{3}$ loops in n dimensions. Since the T_n gate only is defined for prime n the dimensions of greatest interest will be the prime dimensions where $3(n-1)k = n^2 - 1$. Since these seems to be the most information carrying with the fewest number of loops where universality could be implemented. The first few of these perfect-prime dimensions are $n = 5, 11, 17, 23, \dots$. It has already been shown that fault tolerant QC is implementable in odd prime dimensions [25], so a high fidelity controllable scheme in these dimensions would be desirable.

The quantum gate \mathbb{U} can be formulated as a linear combination of the dark state and dark paths just as in the qutrit case,

$$|\psi(t)\rangle = f_0 |d\rangle + \sum_{i=1}^m f_i |D_i(t)\rangle, t \in [0, T] \quad (79)$$

where the coefficients can be solved for by choosing an initial state $|\psi(0)\rangle$. This corresponds to one loop, by using $|\psi(T)\rangle$ as an initial state for the next loop. By iterating this method one can simulate $\mathbb{U} = U^k$ where each loop U have a different set of parameters. So in theory this can be implemented and simulated for any dimensional qudit, just as for the case of the qutrit.

4 Conclusions and outlook

We have shown how to explicitly create a quantum mechanical system which could be used to emulate a qutrit and accompanying single-qutrit gates, this is done by expanding the dark path qubit scheme [15] into a higher dimension. We have shown how it will generalize in the qudit case and using auxiliary states to improve the robustness of the gates. Universality for the qutrit can be obtained using the Hadamard and T -gate equivalent gates in 3 dimensions. The qutrit gates have a high fidelity and their robustness is improved by the inclusion of the auxiliary state in a similar way as for the qubit [15] which suggests that this method can be beneficial for higher dimensional qudits to improve robustness. In the general case for the qudit we have also shown how any dimensional single-qudit diagonal unitary could be created by a single multi-pulse loop in parameter space and that non-diagonal unitaries scale linearly at worst in the number of loops and parameters required for control of each loop scale linearly. The possibility that the scheme expands efficiently into certain prime dimension have been discussed.

One of the greatest challenges of this project was how to choose parameters to replicate the unitary gates. Analytical construction of any diagonal unitary is solved, but there are many other gates that would be interesting to study in more detail. First, finding how different constraints impact the construction of gates and vice versa. Second, how the number of parameter space loops relate to which unitaries are reachable. This problem is very interesting since reduction of loops decreases parameters needed for control, fewer loops implies less error and makes the scheme less costly to run. For both those cases numerical optimization interesting problem to study, will the numerically optimized gates perform worse than analytical ones, and are there any deeper problems associated with this method. Numerical optimization in the high dimensional parameter spaces present in these system are most likely not well-behaved and may contain a lot of false minimums. There might be some structure related to the parametrization that could be used to study the parameter space to gain insight into how to optimally solve the problem.

Another obvious concern that could be addressed is the lack of multi-qudit gates, entanglement is one of the fundamentals of QC and showing a scheme on how these could be constructed would be an obvious improvement. This could maybe achieved through some detuning or some more explicit physical implementation.

Overall many aspects would be interesting to delve deeper into, and I am almost certain that improvements to this scheme would be possible.

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