I. INTRODUCTION

The quantum harmonic oscillator possessing an infinite-dimensional number-state space (i.e., the maximum occupation number $s$ tends to infinity) can well model the Bosonic fields. For example, light is considered a set of infinite number of harmonic oscillators. Since the classical observable phase of light has no corresponding Hermitian operator counterpart (quantum optical phase) [1–3], we will meet with some difficulties when we investigate the number-phase uncertainty relations of the maser and squeezed state in quantum optics. These difficulties may be as follows: (i) the exponential-form operator exp$[i\hat{\phi}]$ is not unitary, where $\hat{\phi}$ is a phase operator; (ii) the number-state expectation value of Dirac’s quantum relation $[\hat{\phi}, \hat{N}] = -i$ is zero, i.e., $\langle n| [\hat{\phi}, \hat{N}] |n \rangle = 0$. Here, $\hat{N}$ denotes the occupation-number operator of photon fields. The expectation value of the right-handed side of Dirac’s quantum relation, i.e., $\langle n|(-i)|n \rangle$ is, however, nonvanishing; (iii) the number-phase uncertainty relation $\Delta N\Delta \phi \geq \frac{1}{2}$ would imply that a well-defined number state would actually have a phase uncertainty of greater than $2\pi$ [4]. In order to obtain a well-behaved Hermitian phase operator, Pegg and Barnett defined a phase state as $|\theta\rangle = \lim_{s \to \infty} (s+1)^{-\frac{1}{2}} \sum_{n=0}^{s} \exp(in\theta)|n\rangle$, where $|n\rangle$ (where $n = 0, 1, 2, ..., s$) are the $s+1$ number states, which span an $(s+1)$-dimensional state space, and thus suggested an alternative, and physically indistinguishable, mathematical model of the monomode field corresponding to a finite but arbitrarily large state space [4]. This means that the state space $\{|n\rangle\}$ with $0 \leq n \leq s$ has a finitely upper level ($|s\rangle$) and the maximum occupation number of particles is $s$ rather than infinity. It was shown that Pegg-Barnett (P-B) approach has several advantages over the conventional Susskind-Glogower formulation [2]. For example, the Pegg-Barnett phase operator is consistent with the vacuum being a state of random phase, while the Susskind-Glogower phase operator does not demonstrate such a property of the vacuum [4]. The resulting number-phase commutator in the P-B approach does not yet lead to any inconsistencies, but satisfies the condition for Poisson-bracket-commutator correspondence. The P-B formulation opened up new opportunities for treating some problems in quantum optics (e.g., atomic coherent population trapping (CPT) and electromagnetically induced transparency (EIT) [5]) and in quantum information [6].

In this paper we will further consider the Pegg-Barnett harmonic oscillator that involves a finite large state space, and show that it possesses an $su(n)$ Lie algebraic structures. Based on this consideration, we will generalize the Pegg-Barnett oscillator to a supersymmetric case.

II. ALGEBRAIC STRUCTURE OF P-B OSCILLATOR

In this section, we will study the algebraic structures of P-B harmonic oscillators with arbitrary occupation number $s$. For the Pegg-Barnett (P-B) harmonic oscillator with a finite but arbitrarily large state space of $s+1$ dimensions, the commutation relation between the annihilation and creation operators is $[a, a^\dagger] = \mathcal{A}$ instead of the conventional commutation relation $([a,a^\dagger] = \mathcal{I}$) of the Bosonic field. This means that the non-semisimple Lie algebra should be generalized to a semisimple one. It can be verified that the matrix representation of the operators $a, a^\dagger$ and $\mathcal{A}$ takes the following form (in the number-state base-vector set)

$$a_{mn} = \sqrt{n}\delta_{m,n-1}, \quad a_m^\dagger = \sqrt{n+1}\delta_{m,n+1},$$

$$\mathcal{A}_{mn} = \delta_{mn} - (s+1)\delta_{m,s}\delta_{n,s},$$

(2.1)

where the subscript $m, n$ (which run from 0 to $s$ only) denote the matrix row-column indices. The remaining generators $\mathcal{M}, \mathcal{M}^\dagger, \mathcal{K}, \mathcal{F}, \mathcal{F}^\dagger, ...$ can be obtained as follows ($0 \leq m, n \leq s$):

$$[a, \mathcal{M}]_{mn} = (s+1)\sqrt{s}(\delta_{m+1,s}\delta_{n,s}) = (s+1)\sqrt{s}\mathcal{M}_{mn},$$

$$[a^\dagger, \mathcal{M}]_{mn} = -(s+1)\sqrt{s}(\delta_{m,s}\delta_{n+1,s}) = -(s+1)\sqrt{s}\mathcal{M}^\dagger_{mn},$$

$$[\mathcal{M}, \mathcal{M}^\dagger]_{mn} = (\delta_{m,s}\delta_{n,s} - \delta_{m+1,s}\delta_{n+1,s}) = -\mathcal{K}_{mn},$$

$$[\mathcal{A}, \mathcal{M}] = (1+s)\mathcal{M}, \quad [\mathcal{A}, \mathcal{M}^\dagger] = -(1+s)\mathcal{M}^\dagger,$$
\[ [a, M]_{mn} = -\sqrt{s-1} \delta_{m+1,s-1} \delta_{ns} = -\sqrt{s-1} F_{mn}, \]
\[ [a^\dagger, M^\dagger]_{mn} = \sqrt{s-1} \delta_{ms} \delta_{n+1,s-1} = \sqrt{s-1} F_{m'n}, \]
\[ [K, F] = -F, \quad [K, F^\dagger] = F^\dagger, \]
\[ [M, K] = 2M, \quad [M^\dagger, K] = -2M^\dagger, \quad \ldots \] (2.2)

In the following, we will take into account the particular cases of \( s = 1, 2 \). According to the commutation relations (2.2) for the case of \( s = 1 \), in the number-state space, the matrix representation of the generators \( a, a^\dagger \) and \( A \) are reduced to \( a = (\sigma_1 + i \sigma_2)/2 \), \( a^\dagger = (\sigma_1 - i \sigma_2)/2 \) and \( A = \sigma_3 \), where \( \sigma_i \)'s (\( i = 1, 2, 3 \)) are Pauli's matrices. In the meanwhile, the rest generators \( M, M^\dagger, K, F, F^\dagger, \ldots \) are \( M = -a, M^\dagger = -a^\dagger, K = -A, F = O, F^\dagger = O, \ldots \), respectively. It is thus shown that the three generators \( \{a, a^\dagger, A\} \) form a Lie algebra. It follows that \( aa^\dagger + a^\dagger a = I \). Since the algebraic generators of su(2) Lie algebra can be constructed in terms of Pauli's matrices, the P-B harmonic oscillator with \( s = 1 \) corresponds to the Fermionic fields and possesses an su(2) Lie algebraic structure.

For the case of \( s = 2 \), it can also be shown that the algebraic generators \( a, a^\dagger, A, M, M^\dagger, K, F, F^\dagger \) form an sl(3) algebra. The eight Gell-Mann matrices, \( \lambda_i \), can therefore be constructed in terms of these algebraic generators, i.e.,
\[ \lambda_1 = a + a^\dagger + \sqrt{2}(M + M^\dagger), \]
\[ \lambda_2 = i[a^\dagger - a + \sqrt{2}(M^\dagger - M)], \]
\[ \lambda_3 = A + 2K, \quad \lambda_4 = F + F^\dagger, \]
\[ \lambda_5 = i(F^\dagger - F), \quad \lambda_6 = -(M + M^\dagger), \]
\[ \lambda_7 = -i(M^\dagger - M), \quad \lambda_8 = \frac{1}{\sqrt{3}} \lambda_8. \] (2.3)

Thus we demonstrated that the P-B harmonic oscillator with \( s = 2 \) possesses an su(3) Lie algebraic structure.

As to the case of arbitrary integer \( s \), it can be verified that the P-B harmonic oscillator possesses an su\((s+1)\) Lie algebraic structure: if \( G \) represents the linear combination of the Hermitian operators, and consequently \( G = \sigma_1^\dagger \), then the exponential-form group element operator \( U = \exp(i G) \) is unitary. Besides, since \( a, a^\dagger \) and \( A \) are traceless, all the generators derived by the commutators in (2.2) and \( G \) are also traceless due to the cyclic invariance in the trace of matrices product. Thus the determinant of the group element \( U \) is unity, i.e., \( \det U = 1 \), because of \( \det U = \exp(i tr(G)) \). Since it is known that such a group element \( U \) that satisfies simultaneously the above two conditions is the group element of the su\((n)\) Lie group, the high-dimensional Gell-Mann matrices, which closes the corresponding su\((n)\) Lie algebraic commutation relations among themselves, can also be constructed in terms of the generators \( a, a^\dagger, A, M, M^\dagger, K, F, F^\dagger, \ldots \) presented above. It is thus concluded that the P-B harmonic oscillator with a maximum occupation number \( s \) has an \((s+1)\)-dimensional number-state space and possesses an su\((s+1)\) Lie algebraic structure.

Considering the case of \( s \to \infty \) is of typically physical interest. Apparently, it is seen that \( \mathcal{A} \) tends to a unit matrix \( I \), and all the remaining generators (except \( a \) and \( a^\dagger \)), the off-diagonal matrix elements of which approach zero, are thus reduced to \( O \). This, therefore, means that the P-B harmonic oscillator with an infinite-dimensional state space just corresponds to the Bosonic fields.

We hope the consideration of algebraic structures of P-B oscillator presented here might be applicable to the investigation of some related topics such as fractional statistics, anyon [7] and cyclic representation of quantum algebra (group) [8].

### III. THE SUPERSYMMETRIC CASE

Here, we will suggest a concept of supersymmetric P-B oscillator. For this aim, we should first take account of a set of algebraic generators \( \{N, N', Q, Q^\dagger, \sigma_3\} \), which possesses a supersymmetric Lie algebraic properties, i.e.,
\[ Q^2 = (Q^\dagger)^2 = 0, \quad [Q, Q^\dagger] = N' \sigma_3, \quad [N, N'] = 0, \]
\[ [N, Q] = -Q, \quad [N, Q^\dagger] = Q^\dagger, \]
\[ \{Q, Q^\dagger\} = N', \quad \{Q, \sigma_3\} = \{Q^\dagger, \sigma_3\} = 0, \]
\[ [N', Q] = 0, \quad [N', Q^\dagger] = 0, \]
\[ [Q, \sigma_3] = -2Q, \quad [Q^\dagger, \sigma_3] = 2Q^\dagger, \]
\[ (Q^\dagger - Q)^2 = -N', \quad [N', \sigma_3] = 0, \] (3.1)
where \( \{ \} \) denotes the anticommuting bracket. One of the physical realization of such a Lie algebra is the multiphoton Jaynes-Cummings model [9,10], the generators of the interaction Hamiltonian of which is \( Q^\dagger = \frac{1}{\sqrt{k}} a^\dagger \sigma_- \) and \( Q = \frac{1}{\sqrt{k}} a k \sigma_+ \), where \( \sigma_\pm = (\sigma_1 \pm i \sigma_2)/2 \), and \( k \) denotes the photon number in each atomic transition process [9–11]. It can be readily verified that the eigenvalue equation of the invariant operator \( N' \) (which commutes with all the operators \( N, Q^\dagger, Q, \sigma_3 \)) is of the form
\[ N' \left( \frac{|m\rangle}{|m+k\rangle} \right) = C_{m+k}^m \left( \frac{|m\rangle}{|m+k\rangle} \right) \] (3.2)
with the eigenvalue \( C_{m+k}^m = \frac{(m+k)!}{m!k!} \). Further calculation shows that
\[ Q^\dagger \left( \frac{|m\rangle}{|m+k\rangle} \right) = \sqrt{\frac{(m+k)!}{m!k!}} \left( \frac{0}{|m+k\rangle} \right), \quad Q \left( \frac{|m\rangle}{|m+k\rangle} \right) = \sqrt{\frac{(m+k)!}{m!k!}} \left( \frac{0}{|m\rangle} \right). \] (3.3)

Thus we obtain the following supersymmetric quasialgebra \( \{N, Q^\dagger, Q, \sigma_3\} \) in a sub-Hilbert-space corresponding to the particular eigenvalue \( C_{m+k}^m \) of the invariant operator \( N' \) by replacing the generator \( N' \) with \( C_{m+k}^m \) in the commutation relations in (3.1), i.e.,
\[ [Q, Q^\dagger] = C^m_{m+k} \sigma_3, \quad \{Q, Q^\dagger\} = C^m_{m+k}, \]
\[
(Q^\dagger - Q)^2 = -C^m_{m+k}. \tag{3.4}
\]

Here we assume that the supersymmetric P-B oscillator can be characterized by the above supersymmetric quasialgebra (3.4). By analogy with the Hamiltonian \( (H = \frac{1}{2} \{a, a^\dagger\} \omega) \) of the Bosonic oscillator, the Hamiltonian of such a supersymmetric P-B oscillator may be written in the form \( H = \frac{1}{2} \{Q, Q^\dagger\} \Omega \), where \( Q^\dagger \) and \( Q \) can be regarded as the creation and annihilation operators and the eigenvalue \( C^m_{m+k} \) of \( \Omega^\dagger \) may be considered the particle occupation number of the supersymmetric P-B oscillator in a certain number state (e.g., the eigenstate of \( \Omega^\dagger \)).

It should be noted that the multiphoton Jaynes-Cummings model, the Hamiltonian of which possesses a supersymmetric Lie algebraic structure, can model the behavior of the supersymmetric P-B oscillator: specifically, by choosing the appropriate parameters for the multiphoton Jaynes-Cummings model, the Hamiltonian of which can take the form of the linear combination of \( Q \) and \( Q^\dagger \), namely, \( H = gQ + g^*Q^\dagger \), where \( g \) and \( g^* \) denote the coupling coefficient of this Jaynes-Cummings model. If the stationary Schrödinger equation of the multiphoton Jaynes-Cummings model is of the form \( \tilde{H}\psi = \epsilon\psi \), then one can easily obtain a new eigenvalue equation \( \{Q, Q^\dagger\}\psi = \epsilon^2/(g^*g)\psi \).

**IV. A POTENTIAL APPLICATION**

In the previous sections, we discussed the Lie algebraic structure of P-B oscillator and then proposed a supersymmetric generalization. Here, we will consider a potential application of the concept of the supersymmetric P-B oscillator to the mass spectrum of charged leptons. First, we suggest a new mass formula for the charged leptons that agrees with experimental values to a high degree of accuracy. This leptonic mass spectrum is constructed based on the following three clues: (i) the supersymmetric P-B oscillator; (ii) Barut’s viewpoint of magnetic self-interaction of charged leptons [12,13]; (iii) the experimental values of charged leptons.

According to the previous section, the energy (or frequency) of the supersymmetric P-B oscillator is proportional to the combination coefficient \( C^m_l \) (i.e., \( m! \sum_{m-l}^{m-l} C^m_l \)). Particularly, if the two integers \( m \) and \( l \) satisfy the relation \( m = l \), then this supersymmetric oscillator will be reduced to a regular two-dimensional Pegg-Barnett oscillator. We assume that there is a deep connection between the supersymmetric Pegg-Barnett oscillator and the generation replication of charged leptons, namely, the supersymmetric Pegg-Barnett oscillator can model some behaviors (at least the mass spectrum) and the internal structures (should such exist) of some certain elementary particles. Thus, the above combination coefficient should be introduced into the leptonic mass spectrum under consideration. In addition, Barut showed that the mass difference between muon and electron may result from the magnetic self-interaction energy of the electron [12,13]. He believed that the radiative effects give an anomalous magnetic moment to the electron, which implies an extra magnetic energy [12,13]. By using the quantization formulation according to the Bohr-Sommerfeld procedure, one can obtain the magnetic energy of a system consisting of both a charge and a magnetic moment as \( E_l = \lambda^2 \) with \( l \) and \( \lambda \) being the angular quantum number of the system and a certain constant coefficient [12,13]. Based on the above enlightening clues, a new mass formula for the charged leptons can be constructed as follows

\[
m_n = \left(1 + \frac{1}{2\alpha} \sum_{l=0}^{n} C^m_l \right) m_e, \tag{4.1}
\]

where \( m_e \), \( \alpha \) and \( n \) denote the electron mass, the electromagnetic fine structure constant and the generation number (generation index of leptons), respectively. Here, the electron (e), muon (\( \mu \)) and tau (\( \tau \)) particle correspond to \( n = 0, 1, 2 \), respectively. By using \( \alpha^{-1} = 137.036 \) and \( m_e = 0.51110 \) MeV, it follows from the formula (4.1) that the masses of charged leptons of various generations are \( m_\mu = 105.55 \) MeV, \( m_\tau = 1786.2 \) MeV and \( m_3 = 4622.2 \) MeV. The experimental values for muon and tau masses are \( m_\mu^{\text{exp}} = 105.66 \) MeV, \( m_\tau^{\text{exp}} = 1784.2 \) MeV [14]. Thus the relative precisions of the formula (4.1) are only of \( -1.04 \times 10^{-3} \) (for muon) and \( +1.12 \times 10^{-3} \) (for tau), respectively. It should be noted that the currently accepted value for the tau lepton mass that was measured in 1992 is \( m_\tau^{\text{exp}} = 1776.9 \) MeV [15]. Someone may therefore argue that the formula (4.1) will not agree very well with such an experimental result. But we will point out that this is not the true case. Since for the case of tau, the interaction energy scale is of GeV, one should consider the running coupling “constant” of electromagnetic interaction (i.e., the variation of the fine structure constant at different energy scale). By taking account of this factor, the inverse of the fine structure constant, \( \alpha^{-1} \), will decrease. As a result, the mass of tau particle obtained by using (4.1) will therefore become less and is still in good agreement with the experimental value obtained in 1992 [15]. In a word, the mass spectrum of charged leptons presented in this paper can agree with the experiments to about one part in \( 10^3 \). It is shown from the mass spectrum (4.1) that there may exist a fourth (and even final) charged lepton, the mass (\( m_3 \)) of which is more than 9000 times that of electron. Even though such a heavy “electron” has so far never been observed experimentally, such a mass spectrum may still be of interest, since the most remarkable feature of (4.1) is that the total generation number of leptons is finite (rather than infinite), which results from
the fact that the integer \( l \) in the combination number \( C_3^l \)
can be taken to be \( l = 0, 1, 2, 3 \) only. So, it might be
able to interpret one of the most fundamental problems
in particle physics and quantum field theory, i.e., the
finite-generation-number phenomenon of fermions. This
advantage has never arisen in the previous mass spec-
tra of leptons proposed by other investigators, where the
mass formulae could not suppress the generation index
and the total number of the allowed generations of the
fermion chain is infinite \([12,13,16]\).

It should be noted that although the present experi-
mental evidences show that the total generation number
of the fermion chain is three \([17]\), a probe into the
potential existence of extra generations of fermions has still
attracted attention of many investigators by now \([18–28]\).
In experiments, the latest electroweak precision data al-
 lows the existence of additional chiral generations in the
standard model \([19]\). Arik \textit{et al.} studied the influence
of extra generations on the production of the standard
model Higgs boson at hadron colliders \([18,19]\). In theo-
etical investigations, some authors extended the standard
electroweak gauge model to include a fourth generation of
fermions, and considered the exotic interactions involv-
ing fourth-generation quarks and leptons which cannot
be confused experimentally with those of the standard
model, or suggested a completely different interaction
model for the extra-generation fermions \([24–33]\). These
studies may provide a possible test of the fourth gen-
eration and would give a signal of new physics.

\section{V. CONCLUSION}

We discuss the \( \text{su}(n) \) Lie algebraic structure of the
P-B oscillator and generalize it to a supersymmetric case.
We think that in some sense the multiphoton Jaynes-
Cummings model can describe the behavior of the su-
persymmetric P-B oscillator. Based on the concept of
supersymmetric P-B oscillator and Barut’s viewpoint of
magnetic self-interaction of charged leptons, we construct
a new mass formula for charged leptons, the most re-
markable feature of which is such that the total number
of the generations of the charged leptons is finite. We
think that since it can present a possible explanation for
the finite-generation phenomenon of charged leptons, our
tentative analysis (the application of supersymmetric P-
B oscillator to leptonic mass spectrum) in the present
paper may still deserve further consideration. We hope
this consideration might provide us with an insight into
the problems such as the physical origin of the generation
replication of the fermion chain.

\textbf{Acknowledgements} This work was supported par-
tially by the National Natural Science Foundation of
China under Project No. 90101024.

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